Nonreciprocity and magnetic-free isolation based on optomechanical interactions

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Nonreciprocal components, such as isolators and circulators, provide highly desirable functionalities for optical circuitry. This motivates the active investigation of mechanisms that break reciprocity, and pose alternatives to magneto-optic effects in on-chip systems. In this work, we use optomechanical interactions to strongly break reciprocity in a compact system. We derive minimal requirements to create nonreciprocity in a wide class of systems that couple two optical modes to a mechanical mode, highlighting the importance of optically biasing the modes at a controlled phase difference. We realize these principles in a silica microtoroid optomechanical resonator and use quantitative heterodyne spectroscopy to demonstrate up to 10 dB optical isolation at telecom wavelengths. We show that nonreciprocal transmission is preserved for nondegenerate modes, and demonstrate nonreciprocal parametric amplification. These results open a route to exploiting various nonreciprocal effects in optomechanical systems in different electromagnetic and mechanical frequency regimes, including optomechanical metamaterials with topologically non-trivial properties.
Lorentz reciprocity stipulates that electromagnetic wave transmission is invariant under a switch of source and observer, and its implications widely permeate physics. To violate reciprocity and obtain asymmetric transmission, suitable forms of time-reversal symmetry breaking are required. In optical and microwave systems this is usually achieved using magneto-optic material responses. However, a vibrant search for alternative methods to break reciprocity, mimicking a magnetic bias, has taken shape in recent years. This is fuelled by the typically weak magneto-optic coefficients in natural materials and/or their associated losses, and the technological promise of integrated on-chip nonreciprocal devices, including isolators and circulators. A promising approach relies on spatiotemporal modulation of the refractive index to break time-reversal symmetry. Such modulation allows imparting a nonreciprocal modulation of the refractive index to break time-reversal symmetry. This allows the definition of minimal conditions to achieve ideal optomechanical nonreciprocity, that meets these minimal conditions are presented, showing the on-chip implementation of an optical isolator and demonstration of a nonreciprocal optomechanical amplifier.

**Results**

**Nonreciprocal mode transfer and optomechanical isolation.** Consider a basic system (Fig. 1a) of two optical modes with frequencies \(\omega_1, \omega_2\), both coupled to a mechanical mode with frequency \(\Omega_m\). The Hamiltonian of this system is

\[
H = \hbar \sum_{j=1,2} \omega_j(x)x_j + \hbar \Omega_m b^\dagger b,
\]

where \(a\) and \(b\) denote the photon and phonon annihilation operators, respectively, and \(\omega_j(x) = \omega_j - \Omega_j\). \(\Omega_j\) is the zero-point motion and \(\Omega_j\) the optical frequency shift per unit displacement. If both optical modes are driven by a strong coherent laser to an intracavity field \(x_j \exp(-i\omega_j t)\), the linearized Hamiltonian in a frame rotating at \(\omega_{\text{control}}\) reads

\[
H_L = -\hbar \sum_{j=1,2} \Delta_j \delta a_j^\dagger \delta a_j + \hbar \Omega_m b^\dagger b
\]

\[
-\hbar \sum_{j=1,2} \left( g_j^d \delta a_j^\dagger b + g_j^c \delta a_j b + g_j^s \delta a_j a_j^\dagger b^\dagger \right).
\]

Figure 1 | Nonreciprocity in a multimode optomechanical system. (a) Optomechanical system: two optical resonators at frequencies \(\omega_1\) and \(\omega_2\), both coupled to a mechanical mode at frequency \(\Omega_m\). (b) General description: the optical modes \((a_1, a_2)\) are coupled to a mechanical mode \(b\) with coupling rates \(g_1\) and \(g_2\). The path \(a_1 \rightarrow b \rightarrow a_2\) picks up a phase \(\Delta \phi = \arg(g_2) - \arg(g_1)\) that is opposite to that of the reversed path \(a_2 \rightarrow b \rightarrow a_1\). Two input/output ports \((s_1, s_2)\) are coupled to the optical modes with rates \(d_{\perp}\), Interfering both paths with direct scattering through the waveguide can build an optical isolator. (c) A ring-resonator supports even and odd optical modes \((a_0, a_2)\), superpositions of clockwise \((a_{+\omega})\) and counter-clockwise \((a_{-\omega})\) propagating modes. As the two modes are \(\pi/2\) out of phase with respect to a wave propagating in the waveguide, a control incident from a single input port fulfills the optimal driving conditions to break reciprocity. The graph sketches the spatial intensity profile of the two modes along the ring of the ring-resonator as a function of the angle \(\theta\) with respect to the dashed line.
excites an equal superposition of even and odd modes with \( \pi/2 \) phase difference, such that the requirement on the control phase to maximally break reciprocity is automatically fulfilled. Note that our choice of the even/odd basis (in contrast to the clockwise/counterclockwise basis considered in other work\(^{10,29,30}\)) immediately reveals the role of a nonreciprocal phase shifter. If a direct pathway exists (Fig. 1c), its interference with the clockwise/counterclockwise basis considered in other systems.

The nature of the nonreciprocal response is determined by the direct coupling between the two channels: if it is forbidden (Fig. 1a), the system primarily functions as a nonreciprocal phase shifter. If a direct pathway exists (Fig. 1c), its interference with the resonant path that collects a nonreciprocal \( \pi \) phase shift enables ideal isolation under appropriate conditions. In our experiment, we demonstrate optical isolation by studying the two-way transmittance of a probe signal at frequency \( \omega_{\text{probe}} \) through a tapered fibre that is coupled to a microcavity (\( \omega_{\text{micro}} = 194.5 \text{ THz} \)) with linewidth \( \kappa/2\pi = 28 \text{ MHz} \). With the control laser incident from one direction, the transmittance is quantified using a heterodyne spectroscopic technique, where a probe beam propagating in the forward or backward direction is recombined with the control, and their beat analysed (see Fig. 2a and Methods). The fact that the measurement technique used here allows to quantify the resulting transmittance provides a means to extract the obtained optical isolation, in contrast to the qualitative measurements reported in\(^{10,30}\).

The resulting probe transmittance (Fig. 2b) for \( \Delta_{1,2} = -\Omega_m \) and near-critical coupling conditions shows a bidirectional transmission dip as the probe frequency is scanned across the cavity resonance. Importantly, the OMIT window\(^{28}\), which results from destructive intracavity interference of anti-Stokes scattering of the control beam from the probe-induced mechanical vibrations with the probe beam itself, is solely present for co-propagating control and probe (dark green circles). For reversed probe direction OMIT is absent (light green squares). The device thus acts as an optical isolator, reaching up to 10 dB of isolation (Fig. 2c).

The nonreciprocal scattering matrix. To predict the magnitude of such nonreciprocal transmission, we use temporal coupled mode theory\(^{39}\) to formulate the scattering matrix \( S \) of a general system described by equation (2), relating input \( (\hat{\delta}_{j+}^\dagger) \) and output \( (\hat{\delta}_{j+}) \) waves at frequency \( \omega_{\text{probe}} \) in the ports \( j = 1, 2 \) via \( \hat{\delta}_{j+}^\dagger \hat{\delta}_{j+}^\dagger = S(\omega_{\text{probe}}) \hat{\delta}_{j+}^\dagger \hat{\delta}_{j+}^\dagger \). The dynamics of a two-mode system described by a linear time-evolution operator

![Diagram of experimental setup and isolation](image_url)
\[ M \text{ reads} \]
\[ \frac{d}{dt} \left( \begin{array}{c} \delta a_1 \\ \delta a_2 \end{array} \right) = iM \left( \begin{array}{c} \delta a_1 \\ \delta a_2 \end{array} \right) + D^T \left( \begin{array}{c} \delta s_1^+ \\ \delta s_2^- \end{array} \right), \]

where \( D \) describes the mutual coupling to the input/output fields. The output fields are found from

\[ \left( \begin{array}{c} \delta s_1^- \\ \delta s_2^+ \end{array} \right) = C \left( \begin{array}{c} \delta s_1^+ \\ \delta s_2^- \end{array} \right) + D \left( \begin{array}{c} \delta a_1 \\ \delta a_2 \end{array} \right), \]

where \( C \) describes the direct coupling between the two ports. Note that these expressions can be related to the quantum optics input/output formalism via a redefinition of the input fields (Supplementary Notes 1, 2 and 6). Here we prescribe the individual optical modes to be reciprocal, such that coupling to in- and outgoing fields is identical\(^\text{39}\). In our system, it necessitates the choice of the even/odd mode basis. In the frequency domain, equations (3 and 4) yield the total scattering matrix

\[ S = C + iD(M + \omega I)^{-1}D^T, \]

with \( I \) the identity matrix, \( M \) the Fourier transform of operator \( M \), and \( \omega = \omega_{\text{probe}} - \omega_{\text{control}} \). In a general two-mode system, the difference between forward and backward complex transmission coefficients thus reads

\[ S_{21} - S_{12} = \frac{i \det(D)(m_{12} - m_{21})}{\det(M + \omega I)}, \]

showing that reciprocity can be broken as long as \( \det(D) \neq 0 \) and \( m_{12} \neq m_{21} \) (with \( m_{ij} \) the elements of \( M \)). This important result identifies the minimal conditions to break reciprocity: a full-rank \( D \) matrix, requiring an asymmetry in the coupling between the two optical modes and the channels \( s_1 \) and \( s_2 \), and an asymmetric evolution matrix, enforcing the coupling rate from mode 1 to mode 2 to be different from that of mode 2 to mode 1. As explained above, this can be implemented through optomechanical interactions.

The evolution matrix \( M \) that describes optomechanical interactions (Fig. 1) is obtained from the equations of motion

\[ \frac{d}{dt} \left( \begin{array}{c} \delta a_1 \\ \delta a_2 \end{array} \right) = i \left( \begin{array}{cc} \tilde{\Delta}_1 + i\kappa_1 / 2 & 0 \\ 0 & \tilde{\Delta}_2 + i\kappa_2 / 2 \end{array} \right) \left( \begin{array}{c} \delta a_1 \\ \delta a_2 \end{array} \right) \right) + i \left( \begin{array}{cc} g_1 (b + b^\dagger) & g_2 (b + b^\dagger) \\ g_2 (b + b^\dagger) & g_1 (b + b^\dagger) \end{array} \right) \left( \begin{array}{c} \delta s_1^+ \\ \delta s_2^- \end{array} \right), \]

\[ \frac{d}{dt} b = (-i\Omega_m - \Gamma_m / 2)b + i\left( g_1 \delta a_1 + g_1 \delta a_1^\dagger + g_2 \delta s_2 + g_2 \delta s_2^\dagger \right) + \sqrt{\Gamma_m b_m}, \]

derived from the linearized Hamiltonian (2) including dissipation and coupling between the mechanical resonator and its thermal bath (rightmost term in equation (8)), which under the experimental conditions studied here can be ignored in the analysis (Methods). Likewise we neglect optical quantum noise. Note that we have set coupling between the optical modes to zero, which can always be realized through diagonalization (see Supplementary Note 4). Solving these equations in the frequency domain, applying the rotating wave approximation and using the input-output relation (4), the evolution matrix \( (M + \omega I) \) for \( \omega \approx \pm \Omega_m \), reads

\[ M + \omega I = \left( \begin{array}{cc} \Sigma_{\omega_0} + \frac{|g_1|^2}{\kappa_1} & \frac{|g_2|^2}{\kappa_2} \\ \frac{|g_2|^2}{\kappa_2} & \Sigma_{\omega_0} + \frac{|g_1|^2}{\kappa_1} \end{array} \right). \]

Here, \( \Sigma_{\omega_0} = \omega + \tilde{\Delta}_1 + i\kappa_1 / 2 \) is the inverse optical susceptibility, \( \Sigma_{\omega_m} = \omega \mp \Omega_m + i\Gamma_m / 2 \) the inverse mechanical susceptibility and \( \Gamma_m \) the mechanical damping rate. Importantly, \( (m_{12} - m_{21}) \propto \sin \Delta \phi \), highlighting the importance of the control phase difference to obtain nonreciprocal transmission.

We define individual cooperativities \( C_j \) by \( C_j \equiv 4|g_j|^2 / (\kappa_j \Gamma_m) \) and the total cooperativity \( C \equiv C_1 + C_2 \). By combining (6) and (9), the asymmetric transmission through a two-mode system

![Figure 3](https://example.com/figure3.png)

**Figure 3 | Power-dependence and nonreciprocal amplification.** (a) Difference between forward/backward transmissivities (measured, red circles and theory \( (|S_{21}|^2 - |S_{12}|^2 \), dashed red line) with respect to cooperativity, directly proportional to the control laser power. Together with an increase in contrast, the insertion losses (blue diamonds) decrease with increasing cooperativity (b) The isolation as a function of probe power sent through the fibre. The physical mechanism behind optical isolation is linear, and thus does not depend on probe power. (c) When the control beam is tuned to the blue side band of the cavity, it can parametrically amplify the probe beam that co-propagates with it through the fibre (dark green circles). In contrast, the counter-propagating probe beam (light green squares) experiences normal cavity extinction, thus yielding a nonreciprocal amplifier. With amplification of \( \sim 3 \) dB, the nonreciprocal difference in transmission is \( \sim 23 \) dB. The solid yellow line is a fit of \( |S_{21}|^2 \) (see Methods) yielding \( (\kappa_{1,2} |g_j|)/2\pi \approx (28 \text{ MHz}, 454 \text{ kHz}) \) and \( \eta_{1,2} \approx 0.46 \).
can be written (Supplementary Notes 3 and 5) as
\[
\delta_{12} = -2 \sin \Delta \phi \sqrt{\eta_1 \eta_2} \times \left( \frac{\delta_+ + \delta_-}{2} \right),
\]
where \( \delta_{12} \equiv (\omega + \Omega_m)/(\gamma_m/2) \) is the fraction of energy mode \( j \) radiates in both output channels. Inspection of equation (10) shows that the magnitude of asymmetric transmission at critical coupling (\( \eta_{1,2} = 1/2 \)) is maximally 1, when the cooperativities are large and equal. These conditions, implemented in our experiment, enable the observed strong optical isolation.

**Figure 4 | Non-degenerate optical modes.** (a) Transmittance of an optical split-mode (splitting \( \sim 8.6 \) MHz) as a function of laser-cavity detuning, obtained by sweeping the laser frequency and measuring the resulting transmittance using an oscilloscope. The horizontal axis is calibrated using the EOM placed in the control arm (see Fig. 2a). The black solid line represents a fit of a double Lorentzian lineshape to the blue data points. The red shading is the area under the two fitted Lorentzian lineshapes. (b) Transmittance of the optical probe beam as a function of control-probe detuning with the control frequency fixed at the blue mechanical side band. When the probe beam co-propagates \((S_{12}^2, \text{dark green circles})\) with the control beam an optomechanically induced absorption window appears, while the oppositely propagating probe \((S_{21}^2, \text{light green squares})\) experiences increased transmission. (c) Asymmetric phase transmission for the same measurement as b. Light green squares correspond to \( \arg(S_{12}) \), the dark green circles to \( \arg(S_{21}) \). The solid lines in b are fitted simultaneously to \( |S_{12}|^2 \) (blue line) and \( |S_{21}|^2 \) (yellow line). The resulting parameters are inserted in \( S_{12} \) (blue line) and \( S_{21} \) (yellow line) to yield the lineshapes in c (Methods).

Dependence on power and detuning and mode degeneracy. For degenerate optical modes and the control field tuned to either mechanical sideband, the maximum contrast between forward and backward transmittance is \( \Delta T = |S_{12}|^2 - |S_{21}|^2 \approx (C^{-1} \pm 1)^2 \) at \( \omega = \pm \Omega_m, \Delta_1 = \pm \Omega_m \), where \( C = 2C_1 = 2C_2 \). The pronounced increase of \( \Delta T \) with increasing \( C \) and concomitant decrease of insertion loss, are confirmed by varying the optical drive power (Fig. 3a). The mechanism has strong potential for near-ideal isolation at negligible insertion losses, for example in optimized silica microtoroids, where \( C \approx 500 \) was demonstrated.\(^2\) Moreover, cooperativity enhances the bandwidth, which is ultimately limited by the optical linewidth.\(^3\) An important aspect of this mechanism is that the isolation is independent of probe power (Fig. 3b), differing fundamentally from mechanisms exploiting static nonlinearity to create asymmetric transmission. Note that noise photons originating from the mechanical thermal bath contribute only 0.4% to the measured probe signal (Methods).

For blue-detuned control \((\Delta_1 = +\Omega_m)\), the probe beam experiences parametric amplification if control and probe are co-propagating, while it is fully dissipated when counter-propagating with the pump, thus yielding a nonreciprocal optical amplifier (Fig. 3c). This feature could pose interesting signal processing functionality, including nonreciprocal narrowband RF filtering and insertion loss compensation.

Importantly, equation (10) shows that strong nonreciprocity can also be obtained without optical degeneracy. If the two modes have different frequency and/or linewidth, an optimal control frequency can be chosen to satisfy \( \delta_1 = -\delta_2 = \beta \). Then asymmetric transmittance is maximally
\[
\Delta T = \frac{C(C \pm 2\beta^2)}{(1 \pm C + \beta^2)^2},
\]
showing that larger cooperativity can compensate the effects of mode splitting for \( \beta > 1 \). Figure 4 shows nonreciprocal amplitude and phase transmission with a split optical mode. A probe beam tuned between the even and odd mode frequencies excites both modes with unequal phases. These opposing phases are added to the \( a_1 \rightarrow a_2 \) and \( a_2 \rightarrow a_1 \) optomechanical mode conversion processes, respectively, changing the interference condition with the nonresonant transmission. As a result, both co- and counter-propagating probe fields now interact with the mechanical mode. For a blue-detuned control beam, this yields induced absorption for the co-propagating probe and induced transparency for the counter-propagating probe (Fig. 4b). Note that the induced absorption for the co-propagating beam is related to the relatively low coupling rate (\( \eta_{1,2} < 0.5 \)). It can be turned into gain, as presented in Fig. 3c, for \( \eta_{1,2} > 0.5 \) and/or for increased optical control power. Crucially, since for our system the relation \( \beta \ll \Omega_m/\kappa_{1,2} \) holds (Methods), the deviation of \( \Delta \phi \) from optimal is only 0.2%. As such, a control beam incident from one side still ensures \( \Delta \phi \approx \pi/2 \) and \( C_1 \approx C_2 \), thus fulfilling the requirements for optimal nonreciprocity and maximizing the contrast between forward and backward transmission. In a more general case, optimal conditions may be implemented, for example by supplying control fields with suitable phase and amplitude through both input waveguides. Importantly, the fact that nonreciprocity can be obtained without optical degeneracy increases the range of systems that may be employed.

**Discussion**

We stress that the demonstrated principles are not limited to the experimental implementation using ring resonators shown here, but can be realized in a wide range of optomechanical platforms\(^4\) such as LC circuits\(^5\) and photonic crystal resonators\(^1\). In fact, the high (GHz) frequency of such devices has the prospect of enhancing the bandwidth with respect to the relatively narrow range demonstrated here, towards a range commensurate with typical signal modulation rates. While the resonant nature of the demonstrated mechanism is of course a limit to the general application capability, we foresee several applications that could benefit from magnetic-free isolation over a finite bandwidth. These include in particular...
the protection of on-chip monochromatic laser sources and, with
ground-state cooling\textsuperscript{31,32} or in the strong coupling regime\textsuperscript{23},
low-loss routing of signals carrying quantum information at
negligible added noise, either at optical or microwave
frequencies\textsuperscript{13,29}.

The specific nonreciprocal functionality is governed by the way
these systems are coupled to input/output channels. This, in turn,
is directly related to the nonresonant scattering matrix $C$, as
reciprocity of optical modes dictates $CD^* = -D$ (ref. 39). For the
scenarios in Fig. 5a, described by a diagonal $C$ matrix, each
waveguide couples to a single optical mode, and the system
operates as a nonreciprocal phase shifter (gyrator) in the absence
of other optical loss. Importantly, an on-chip gyrator that is
placed in one arm of an integrated Mach-Zehnder interferometer
could be used to build an on-chip circulator\textsuperscript{42}. In contrast,
we note that nonreciprocity occurs also outside the resolved
side-band regime, although the behaviour there is more complex
due to mixing of sidebands at $\pm \omega$.

In conclusion, we demonstrated and quantified nonreciprocal
transmission through a compact optomechanical isolator and
parametric amplifier, and developed a general theory explaining
the mechanism and unifying the description of various
implementations of optomechanical nonreciprocity in multimode
systems. Our findings identify two general requirements for
any optomechanical system to optimally break reciprocity:
asymmetric coupling of the optical modes to input/output
channels, and a drive phase-difference of $\pi/2$. Since the
requirements for optimal nonreciprocity derived here do not
rely on the handedness of optical\textsuperscript{29,30} or mechanical\textsuperscript{10,11} modes,
our theoretical formalism can be used to realize optomechanical
nonreciprocity in systems that do not exhibit circular symmetry
(Fig. 5). Extending the demonstrated principles to more modes or
channels would enable a variety of nonreciprocal functionality for
both light and sound, including on-chip circulation, gyration\textsuperscript{31}
and enhanced isolation bandwidth. Finally, these nonreciprocal
systems can form the unit cell of optomechanical metamaterials
with topologically non-trivial properties, where the nonreciprocal
phase takes the role of an effective gauge field to establish new
phases for sound and light\textsuperscript{33,34}.

Note added in proof: After submission, we became aware of
related work by Fang et al.\textsuperscript{43} that reports nonreciprocal
transmission in an optomechanical crystal circuit that relies on
the same principle with blue-detuned control.

Methods

Coupling matrix and drive condition in ring resonator. Time-reversal symmetry
and energy conservation dictate that $CD^* = -D$ and $D^*D = \text{diag}(g_{\ell}k_{1\ell}, g_{\ell}k_{2\ell})$\textsuperscript{38}.
Applying these to the even and odd optical modes ($\delta a_1, \delta a_2$) of an evanescently
coupled ring resonator, and choosing $\epsilon_{12} = \epsilon_{21} = 1$, constrains the coupling
matrix $D$ to

$$D = \frac{1}{\sqrt{2}} \begin{pmatrix} i\sqrt{\eta_{k_{11}}} & -i\sqrt{\eta_{k_{22}}} \\ i\sqrt{\eta_{k_{11}}} & i\sqrt{\eta_{k_{22}}} \end{pmatrix}. \tag{12}$$

Together with equations (5 and 9), this $D$ matrix yields the complete expressions for
the scattering matrix elements

$$S_{j\ell} = \epsilon_j + \frac{2A_{\ell}}{\sqrt{g_{\ell}k_{j\ell}}} \left(d_{\ell}d_\ell^* e^{i\delta \ell} + d_{\ell}d_\ell^* e^{-i\delta \ell} \right), \tag{13}$$

where $A_{\ell}$ is given by

$$A_\ell = d_{\ell}d_\ell^* \sqrt{\frac{g_{\ell}}{k_{j\ell}}} (\delta + i(\delta + i) = C_j + d_{\ell}d_\ell^* \sqrt{\frac{g_{\ell}}{k_{j\ell}}} (\delta + i(\delta + i) = C_j, \tag{14}$$

used to fit the data in Figs 2–4.

For a single drive field with amplitude $\tilde{\xi}_\text{control}$ incident through port 1 and using
$G_1 = G_2 = G$, the coupling rates $g_1$ and $g_2$ are given by

$$\begin{pmatrix} g_1 \\ g_2 \end{pmatrix} = G_{\text{opt}} \begin{pmatrix} \xi_1 \\ \xi_2 \end{pmatrix},$$

$$= G_{\text{opt}} \begin{pmatrix} \Sigma_{\ell=0}^{\ell=1} (\omega = 0) \\ 0 \end{pmatrix} D^\dagger \begin{pmatrix} \tilde{\xi}_\text{control} \\ 0 \end{pmatrix} \tag{15}$$

Thus for large detuning $|A_{1,2}| \gg k_{1,2}$, the optimal phase difference $\Delta \phi = \pi/2$ is
automatically satisfied by pumping through a single channel.

Experimental set-up. The silica microtoroid (diameter 41 $\mu$m) is fabricated
using techniques as previously reported (see for example ref. 37). A tuneable
fibre-coupled external cavity diode laser (New Focus, TLB-6728) is locked (using
the electro-optic modulators) to a mechanical sideband of a whispering gallery
mode at 1.542 nm using a modified Pound-Drever-Hall scheme that can be used
independent of the probe beam direction. The probe light is generated using a
commercial double-parallel Mach-Zehnder interferometer (Thorlabs, LN86S-FC)
operated in single-side-band carrier-suppressed mode, driven by the output of a vector network analyser (VNA) at frequency \(\Omega\). The resulting probe light has frequency \(\Omega_{\text{probe}} = \Omega + \Delta\). The sign of the frequency shift \(\Delta\), as well as the suppression of the carrier (by 50 dB with respect to the generated probe) is controlled by bias voltages applied to the double-parallel Mach–Zehnder interferometer. Pump and probe amplitude and polarization are controlled with variable optical attenuators and fibre polarization controllers (Fig. 2a). The probe beam propagating in forward or backward direction is recombined with the control beam and their beat on fast (125 MHz) low-noise photo receivers (D1/D2) is analysed with a VNA. It should be noted that fluctuations of the optical length difference of probe and control paths generate phase fluctuations of the beat analysed by the VNA. To minimize these phase fluctuations on the time scale of the inverse bandwidth (5 kHz) \(^{-1}\) of the VNA, the lengths of the paths Laser/C1/D2 and Laser/0/Photonics/0/Non-reciprocity/0/Photonic/0/Aharonov-Bohm effect based on dynamic modulation. Phys. Rev. Lett. 108, 153901 (2012).

Estep, N. A., Souzas, D. L., Soric, J. & Alù, A. Magnetic-free non-reciprocity and isolation based on parametrically modulated coupled-resonator loops. Nat. Phys. 10, 923–927 (2014).

Souzas, D. L. & Alù, A. Angular-momentum-based nanorings to realize magnetic-free integrated optical isolation. ACS Photonics 1, 198–204 (2014).

Carmon, T., Rokhsari, H., Yang, L., Kippenberg, T. J. & Vahala, K. J. Temporal behavior of radiation-pressure-induced vibrations of an optical microcavity phonon mode. Phys. Rev. Lett. 94, 223902 (2005).

Kippenberg, T. J., Rokhsari, H. Carmon, T., Scherer, A. & Vahala, K. J. Analysis of radiation-pressure induced mechanical oscillation of an optical microcavity. Phys. Rev. Lett. 95, 033901 (2005).

Hafezi, M., Kippenberg, T. J. & Marquardt, F. Cavity optomechanics. Rev. Mod. Phys. 86, 1391–1454 (2014).

Teufel, J. D. et al. Sideband cooling of micromechanical motion to the quantum ground state. Nature 475, 359–363 (2011).

Chan, J. et al. Laser cooling of a nanomechanical oscillator into its quantum ground state. Nature 478, 89–92 (2011).

Vayssagne, E., Deléglise, S., Weihs, S., Schliesser, A. & Kippenberg, T. J. Quantum-coherent coupling of a mechanical oscillator to an optical cavity mode. Nature 482, 63–67 (2012).

Massel, F. et al. Microwave amplification with nanomechanical resonators. Nature 480, 351–354 (2011).

Dong, C. et al. Quantum-coherent coupling of a mechanical oscillator to an optical cavity mode. Nature 482, 63–67 (2012).

Lecoq, F., Clark, J. B., Simmonds, R. W., A. J. Teufel, J. D. Mechanically mediated microwave frequency conversion in the quantum regime. Phys. Rev. Lett. 116, 043601 (2016).

Weis, S. et al. Optomechanically induced transparency. Science 330, 1520–1523 (2010).

Hafezi, M. & Rabl, P. Optomechanically induced non-reciprocity in microring resonators. Opt. Express 20, 7672–7684 (2012).

Shen, Z. et al. Experimental realization of optomechanically induced non-reciprocity. Nat. Photon. 10, 657–661 (2016).

Habraken, S. J. M., Stannigel, K., Lukin, M. D., Zoller, P. & Rabl, P. Continuous mode cooling and phonon routers for phononic quantum networks. New J. Phys. 14, 115004 (2012).

Xu, X.-W. & Li, Y. Optical nonreciprocity and optomechanical circulator in three-mode optomechanical systems. Phys. Rev. A 91, 053854 (2015).

Peano, V., Brendel, C., Schmidt, M. & Marquardt, F. Topological phases of sound and light. Phys. Rev. X 5, 031011 (2015).

Schmidt, M., Kessler, S., Peano, V., Painter, O. & Marquardt, F. Optomechanical creation of magnetic fields for photons on a lattice. Optica 2, 635–641 (2015).

Grudinin, I. S., Lee, H., Hwang, D., Vahala, K. J. Photon laser action in a tunable two-level system. Phys. Rev. Lett. 104, 083901 (2010).

Armani, D. K., Kippenberg, T. J., Spillane, S. M. & Vahala, K. J. Ultra-high-Q toroid microcavity on a chip. Nature 421, 925–928 (2003).

Schliesser, A., Rivière, R., Anetsberger, G., Arcizet, O. & Kippenberg, T. J. Resolved-sideband cooling of a micromechanical oscillator. Nat. Phys. 4, 415–419 (2008).

Fang, K. & Fan, S. Optical isolation based on nonreciprocal phase shift induced by interband photon transitions. Nat. Photon. 3, 91–94 (2009).

Suh, W., Wang, Z. & Fan, S. Temporal coupled-mode theory and the presence of non-orthogonal modes in lossless multimode cavities. IEEE J. Quant. Electron 40, 1511–1518 (2004).

Manipatruni, S., Robinson, J. T. & Lipson, M. Optical nonreciprocity in optomechanical structures. Nat. Nanotechnol. 3, 213-216 (2008).

Fan, L. et al. An All-silicon passive optical diode. Science 335, 447–450 (2012).

Hogan, C. L. The ferromagnetic Faraday effect at microwave frequencies and its applications: the microwave gyror. Bell Syst. Tech. J. 31, 1–31 (1952).

References

1. Deak, L. & Fulop, T. Reciprocity in quantum, electromagnetic and other wave processes. Rev. Lett. 117, 083901 (2016).

2. Haldane, F. D. M. & Rugh, S. Possible realization of directional optical waveguides in photonic crystals with broken time-reversal symmetry. Phys. Rev. Lett. 100, 013904 (2008).

3. Poulton, C. G. et al. Design for broadband on-chip isolator using stimulated Brillouin scattering in dispersion-engineered chalcogenide waveguides. Opt. Express 20, 21235–21246 (2012).

4. Souzas, D. L., Caloz, C. & Alù, A. Giant non-reciprocity at the subwavelength scale using angular momentum-based metamaterials. Nat. Commun. 4, 2407 (2013).

5. Li, E., Eggleton, B. J., Fang, K. & Fan, S. Photonic Aharonov-Bohm effect in photonic phonon interferometers. Nat. Commun. 5, 3225 (2014).

6. Tzeng, L. D., Fang, K., Nussenzveig, P., Fan, S. & Lipson, M. Non-reciprocal phase shift induced by an effective magnetic flux for light. Nat. Photon. 8, 701–705 (2014).

7. Siwa, K. M. et al. Reconfigurable Josephson circulator/directional amplifier. Nat. Commun. 4, 2407 (2013).

8. Guo, X., Zou, C.-L., Jung, H. & Tang, H. X. On-chip strong coupling and efficient frequency conversion between telecom and visible optical modes. Phys. Rev. Lett. 117, 123902 (2016).
43. Fang, K. et al. Generalized nonreciprocity in an optomechanical circuit via synthetic magnetism and reservoir engineering. Preprint at http://arxiv.org/abs/1608.03620 (2016).

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Author contributions

F.R. fabricated the samples and carried out the experiments. F.R. and E.V. developed the experimental set-up. M.-A.M. developed the theoretical model, with contributions from E.V., A.A. and F.R. F.R. and M.A.M. analysed the data. E.V. and A.A. supervised the project. All authors contributed to the writing of the manuscript.

Additional information

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