Loss-induced nonreciprocity

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Abstract
Nonreciprocity is important in both optical information processing and topological photonics studies. Conventional principles for realizing nonreciprocity rely on magnetic fields, spatiotemporal modulation, or nonlinearity. Here we propose a generic principle for generating nonreciprocity by taking advantage of energy loss, which is usually regarded as harmful. The loss in a resonance mode induces a phase lag, which is independent of the energy transmission direction. When multichannel lossy resonance modes are combined, the resulting interference gives rise to nonreciprocity, with different coupling strengths for the forward and backward directions, and unidirectional energy transmission. This study opens a new avenue for the design of nonreciprocal devices without stringent requirements.

Introduction
Optical nonreciprocity, which prohibits a light field from returning along its original path after passing through an optical system in one direction, implying the breaking of the Lorentz reciprocity theorem, is crucially important for both fundamental studies and applied sciences1–3. For example, nonreciprocal devices, such as optical isolators1, optical circulators4, and directional amplifiers5, play important roles in optical communication and optical information processing. Moreover, the topological properties exhibited by nonreciprocal devices make them promising platforms for studying topological photonics6,7 and chiral quantum optics8. To date, a number of approaches have been suggested for generating nonreciprocity, including the use of parity-time (PT)-symmetric nonlinear cavities9,10, spinning resonators11,12, optomechanical interactions13–20, cavity magnonic interactions21, effective gauge fields22,23, and the thermal motion of hot atoms24,25. Despite these achievements, the basic principles for realizing optical nonreciprocity remain limited as a result of the time-reversal symmetry and linear nature of Maxwell’s equations. The existing approaches can be grouped into three categories with the following requirements1–3: (i) magnetic-field-induced breaking of time-reversal symmetry26–29, (ii) spatiotemporal modulation of system permittivity30–36, and (iii) nonlinearity37,38. However, these principles either encounter difficulties in integration39, require stringent experimental conditions40, or have limited performance38. Therefore, it is crucial to break the Lorentz reciprocity theorem by going beyond these approaches.

Here we devise a new principle for realizing optical nonreciprocity by making use of loss. Although it is obvious that loss breaks time-reversal symmetry, it is generally believed that loss cannot lead to optical nonreciprocity as a result of restricted time-reversal symmetry2,3, in which the field amplitudes are reduced while the field ratios are conserved. However, we show that loss under multiple channels with interference gives rise to optical nonreciprocity. The basic principle is that the phase lag induced by loss, which is independent of the energy propagation direction, results in different interference outcomes for the forward and backward directions. In our scheme, neither a magnetic field, the spatiotemporal modulation of permittivity nor nonlinearity is required. On the contrary, the resource we take advantage of is simply energy loss, which is regarded as harmful and undesirable in most studies. This is also different from PT-symmetric schemes in which the
nonreciprocity originates from nonlinear gain with saturation. Our scheme is universal for a variety of physical systems, such as optical cavities and waveguides. This study paves the way for the observation of nonreciprocity and corresponding device design in lossy systems without stringent conditions, and provides opportunities for studying chiral and topological properties in systems with lossy coupling.

Results

As illustrated in Fig. 1a, we consider a generic system in which an array of main resonance modes $a_m$ ($m = 1, 2, \ldots, M$) are linked by a series of connecting modes $c^{(n)}_m$ ($n = 1, 2, \ldots, N$) with decay rates $\kappa^{(n)}_m$. This model can be implemented in a variety of systems, such as optical cavities\textsuperscript{41}, superconducting circuits\textsuperscript{42}, mechanical resonators\textsuperscript{43}, and atomic ensembles\textsuperscript{44}. In the frame rotating at the input laser frequency $\omega_1$, the system Hamiltonian is given by ($\hbar = 1$)

$$
H = -\sum_{m=1}^{M} \Delta_m a_m^\dagger a_m - \sum_{m=1}^{M-1} \sum_{n=1}^{N} \Delta_m^{(n)} c_m^{(n)} c_{m+1}^{(n)} + \sum_{m=1}^{M-1} \sum_{n=1}^{N} \left[ \sum_{l=1}^{N} \left( g_{L,m}^{(n)} a_m^\dagger b_{l,m}^{(n)} + g_{R,m}^{(n)} c_{m+1}^{(n)} c_{l,m}^{(n)} \right) \right] + H.c.
$$

(1)

where $\Delta_m \equiv \omega_1 - \omega_m$ and $\Delta_m^{(n)} \equiv \omega_n - \omega_m^{(n)}$ represent the detunings, with $\omega_m$ ($\omega_m^{(n)}$) being the resonance frequency of mode $a_m$ ($c_m^{(n)}$) and $g_{L,m}^{(n)}$ ($g_{R,m}^{(n)}$) being the coupling coefficient between $a_m$ and $c_{m+1}^{(n)}$. Here, the indices of the main (connecting) modes are denoted by subscripts (superscripts in parentheses) to avoid confusion.

In the above model, $N$ connecting modes are required to realize $N$-channel coupling between $a_m$ and $a_{m+1}$. As sketched in Fig. 1b, this multichannel coupling can also be realized by using only one connecting mode $c_m$ with synthetic frequency dimensions, where $N$ pairs of coupling detunings $\delta_{L,R,m}$ play the role of $N$ coupling channels. In this case, the system is expressed by

$$
H = -\sum_{m=1}^{M} \Delta_m c_m^\dagger c_m - \sum_{m=1}^{M-1} \sum_{n=1}^{N} \Delta_m^{(n)} c_m^{(n)} c_{m+1}^{(n)} + \sum_{m=1}^{M-1} \sum_{n=1}^{N} \left[ \sum_{l=1}^{N} \left( g_{L,m}^{(n)} e^{i\delta_{L,m}^{(n)}a_m^\dagger b_{l,m}^{(n)}} + g_{R,m}^{(n)} e^{-i\delta_{R,m}^{(n)}c_{m+1}^{(n)} c_{l,m}^{(n)}} \right) \right] + H.c.
$$

where $\Delta_m \equiv \omega_1 - \omega_m$ is the detuning of the connecting mode $c_m$. By expressing the connecting mode as $c_m = \sum_{n=1}^{N} \left( c_{L,m}^{(n)} e^{-i\delta_{L,m}^{(n)}} + c_{R,m}^{(n)} e^{i\delta_{R,m}^{(n)}} \right)$, where the $c_{L,R,m}^{(n)}$ are the components corresponding to the coupling detunings $\delta_{L,R,m}^{(n)}$, this Hamiltonian can ultimately be reduced to Eq. (1) under the rotating-wave approximation (see the Supplementary Information for details).

When the detunings $\Delta_m^{(n)}$ or decay rates $\kappa_m^{(n)}$ of the connecting modes are much larger than the coupling rates ($|\Delta_m^{(n)} + i\kappa_m^{(n)} / 2| >> |g_{L,R,m}^{(n)}|$), the connecting modes $c_m^{(n)}$ can be adiabatically eliminated\textsuperscript{45,46}, leading to an effective non-Hermitian Hamiltonian $H_{\text{eff}} = H_{\text{eff}}^0 + H_{\text{eff}}^\text{int}$ (see the Supplementary Information for details). Here,

$$
H_{\text{eff}}^0 = -\sum_{m=1}^{M} (\Delta_m + i\gamma_m / 2 + \Omega_m) a_m^\dagger a_m
$$

is the free-energy term, with $\gamma_m$ being the intrinsic linewidth of mode $a_m$ and $\Omega_m = -\sum_{n=1}^{N} \left[ |g_{L,m}^{(n)}|^2 / (\Delta_m^{(n)} + i\kappa_m^{(n)} / 2) + |g_{R,m}^{(n)}|^2 / (\Delta_m^{(n)} + i\kappa_{m-1}^{(n)} / 2) \right]$ being the resonance shift and broadening; $H_{\text{eff}}^\text{int}$ is the interaction term, given by

$$
H_{\text{eff}}^\text{int} = \sum_{m=1}^{M-1} \left( h_{m+1,m} a_m^\dagger a_m + h_{m+1,m} a_m^\dagger a_{m+1} + h_{m+1,m} a_m^\dagger a_{m+1} \right)
$$

(2)

where $h_{m+1,m}$ ($h_{m+1,m}$) is the effective coupling coefficient for the forward (backward) direction between $a_m$ and $a_{m+1}$. It is seen that the total effective coupling coefficient is a sum of the effective coupling coefficients for each coupling channel. For the $n$-th channel, the amplitude of the effective coupling coefficient is

$$
G_n^{(m)} \equiv |g_{L,m}^{(n)} g_{R,m}^{(n)}| / \sqrt{\Delta_m^{(n)} + \kappa_{m-1}^{(n)} / 4},
$$

where the phase factors include two components: $\phi_m^{(n)}$ and $\theta_m^{(n)}$. The first component, $\phi_m^{(n)} \equiv \arg \left( g_{L,m}^{(n)} g_{R,m}^{(n)} \right)$, refers to the coherent coupling phase, which changes its sign when the coupling direction is reversed as a result of energy conservation. The second component, $\theta_m^{(n)} \equiv \arg (\Delta_m^{(n)} + i\kappa_{m-1}^{(n)} / 2)$, represents the phase lag induced by loss (the loss phase). It is noteworthy that this loss phase is determined only by the loss-detuning ratio, not by the coupling direction, as the losses play the same role for both the forward and backward couplings.

In the absence of losses, i.e., when $\kappa_m^{(n)} = 0$ and thus $\theta_m^{(n)} = 0$, $h_{m+1,m}$ and $h_{m+1,m}$ are complex conjugates, and $H_{\text{eff}}^\text{int}$ is Hermitian. In the presence of losses, the phase lag $\theta_m^{(n)}$ yields a non-Hermitian $H_{\text{eff}}^\text{int}$ with $h_{m+1,m}$ $= h_{m+1,m}$, meaning that a nonreciprocal energy flow does not exist. This is consistent with the conclusion in the previous literature\textsuperscript{2,3}, in which lossy systems are regarded as reciprocal in terms of restricted time-reversal symmetry. However,
when more than one channel exists, the interference between different channels ultimately leads to unequal effective coupling amplitudes for the forward and backward couplings, i.e., $|h_{m,m+1}| \neq |h_{m+1,m}|$ for $N \geq 2$. As depicted in Fig. 1c, the existence of the loss phase ensures that the interference properties are different for the forward and backward couplings. By tuning the phases, the forward coupling can be made to experience constructive interference, whereas the backward coupling will undergo destructive interference, leading to a nonzero forward coupling strength but a backward coupling strength of zero. This asymmetric coupling leads to an asymmetric scattering matrix (see the Supplementary Information for details), which indicates that the Lorentz reciprocity is broken\(^1\). Thus, nonreciprocity can be realized in a lossy system with multichannel interference. It is noteworthy that neither a magnetic field, the spatiotemporal modulation of permittivity, nor nonlinearity is required, and the nonreciprocity originates purely from the losses, which break the time-reversal symmetry.

Without loss of generality, in the following, we consider a two-channel situation ($N = 2$). To realize complete nonreciprocity, the amplitudes of the coupling coefficients for the two channels should be the same, i.e., $G_m^{(1)} = G_m^{(2)}$ (denoted by $G$ in the following), so that complete destructive interference can be achieved. In this case, the forward and backward coupling coefficients are given by

$$h_{\pm} = 2Ge^{\mp i\varphi_0} \cos \frac{\Delta \phi \pm \Delta \theta}{2}$$  \hspace{1cm} (3)$$

where $h_-$ ($h_+$) is short for $h_{m+1,m}$ ($h_{m,m+1}$), $\Delta \phi = \phi_m^{(2)} - \phi_m^{(1)}$ is the difference between the coherent coupling phases for the two channels, $\Delta \theta = \theta_m^{(2)} - \theta_m^{(1)}$ is the loss phase difference, and $\bar{\varphi} = (\phi_m^{(1)} + \phi_m^{(2)})/2$ and $\bar{\theta} = (\theta_m^{(1)} + \theta_m^{(2)})/2$ are the corresponding average phases. Here, the subscripts “$m$” for the $m$-th main mode are omitted for convenience, as we mainly focus on the $M = 2$ case. It is clearly revealed that the phase difference leads to distinct interference patterns for the forward and backward couplings. The conditions for unidirectional nonreciprocal coupling are given by

$$\Delta \phi \mp \Delta \theta = \pi + 2k\pi \hspace{1cm} \Delta \varphi = p\pi, \Delta \theta = q\pi$$  \hspace{1cm} (4)$$

where $k$, $p$, and $q$ are integers. Here, “-” corresponds to unidirectional forward coupling with $|h_-| \neq 0$ and
$|h_-|=0$, whereas “+” corresponds to unidirectional backward coupling with $|h_-|=0$ and $|h_+|\neq 0$.

In Fig. 2a, we plot the amplitudes of the forward and backward couplings $|h_{\pm}|$ as functions of $\Delta \phi$ for various loss phase differences $\Delta \theta$. When $\Delta \theta=0$ or $\pi$, the curves of $|h_{\pm}|$ coincide and thus, nonreciprocity is not achievable. However, for other values of $\Delta \theta$, the difference between $|h_-|$ and $|h_+|$ becomes significant. The curves of $|h_{\pm}|$ have the same lineshape, but the positions are shifted to the left (right) by $\Delta \theta$, showing that nonreciprocity can be realized. As depicted in Fig. 2b, the nonreciprocity ratio $|h_-|/|h_+|$ reaches a maximum/minimum when $\Delta \phi=\pm (\Delta \theta - \pi)$, in agreement with Eq. (4). For example, $\Delta \theta=\pi/2$ and $\Delta \phi=-\pi/2$ lead to unidirectional forward (backward) coupling (third row of Fig. 2). The nonreciprocal properties can also be illustrated by the trajectories in the parameter space expanded by $h_{\pm}$ as functions of $\Delta \phi$. The parameters are $\Delta \theta=0$ (first row), $\Delta \theta=\pi/4$ (second row), $\Delta \theta=\pi/2$ (third row), $\Delta \theta=3\pi/4$ (fourth row), and $\Delta \theta=\pi$ (fifth row).

Fig. 2 Forward and backward coupling strengths. a Amplitudes of the forward coupling $|h_+|$ (red solid curve) and the backward coupling $|h_-|$ (blue dashed curve) as functions of $\Delta \phi$. b Nonreciprocity ratio $|h_-|/|h_+|$ vs. $\Delta \phi$. c Parametric plots of $h_-$ and $h_+$ as functions of $\Delta \phi$. The parameters are $\Delta \theta=0$ (first row), $\Delta \theta=\pi/4$ (second row), $\Delta \theta=\pi/2$ (third row), $\Delta \theta=3\pi/4$ (fourth row), and $\Delta \theta=\pi$ (fifth row).
The unidirectional coupling strength can strongly determine the unidirectional coupling strength. Thus, the optimal condition can be re-expressed as

$$|h_{-}| = 2G|\sin \Delta \theta|$$

As illustrated in Fig. 3a, the loss phase difference \(\Delta \theta\) strongly determines the unidirectional coupling strength. The best performance is achievable for \(\Delta \theta = \pi/2\), with the corresponding maximum being \(|h_{-}|_{\text{max}} = 2G\). It is noteworthy that the unidirectional coupling strength can remain large over a broad range of \(\Delta \theta\), with the full width at half maximum being \(2\pi/3\) (shaded region). By expressing \(\Delta \theta\) in terms of the detunings and losses for the two coupling channels (Fig. 3b), we obtain

$$\tan \Delta \theta = 2\left[ (\Delta^{(1)}k^{(2)} - \Delta^{(2)}k^{(1)}) / [4\Delta^{(1)}\Delta^{(2)} + k^{(1)}k^{(2)}] \right].$$

Thus, the optimal condition can be re-expressed as

$$\frac{\Delta^{(1)}}{k^{(1)}} = -\frac{\Delta^{(2)}}{k^{(2)}} \quad \frac{\Delta^{(1)}}{k^{(1)}} = -\frac{\Delta^{(2)}}{k^{(2)}}$$

whereas the condition for vanishing nonreciprocity is \(\Delta^{(1)}k^{(2)} = \Delta^{(2)}k^{(1)}\). The contour map of \(|h_{-}|\) as a function of \(\Delta^{(1)}/k^{(1)}\) and \(\Delta^{(2)}/k^{(2)}\) is plotted in Fig. 3c, and typical curves for \(|h_{-}|\) as a function of \(\Delta^{(2)}/k^{(2)}\) with a fixed \(\Delta^{(1)}/k^{(1)}\) are plotted in Fig. 3d. These figures show that the parameter ranges for achieving a large \(|h_{-}|\) are very broad. For a pure lossy system, two detunings with opposite signs are preferred, i.e., in one channel, the connecting mode is red detuned, whereas in the other channel the connecting mode should be blue detuned.

The nonreciprocity can also be reflected in the eigenvalues and eigensystems of the modes. In Fig. 4, we plot the real and imaginary parts of the energy eigenvalues and the expansion coefficients of the eigenvectors for the effective Hamiltonian (solid curves) and the original Hamiltonian (dotted curves) as functions of \(\Delta \phi\). The red solid and magenta dotted curves illustrate the eigenvalues for each eigenmode (\(\alpha_m, \beta_m\)), and the blue solid and cyan dotted curves show the distribution in mode \(a_{m+1}\) for each eigenmode (\(\beta_m\)). The parameters are \(k^{(1)} = k^{(2)} = k\), \(\gamma_n = \gamma_{n+1} = 10^{-3}\), \(\Delta^{(1)}/k = 80\), \(\Delta^{(2)}/k = -k/[4\Delta^{(1)}\Delta^{(2)} + k^{(1)}k^{(2)}]\), \(g^{(1)}/k = g^{(2)}/k = 0.1\sqrt{\Delta^{(1)}/k + i/2}\), and \(g^{(1)}/k = g^{(2)}/k = 0.1\sqrt{\Delta^{(2)}/k + i/2}\). This is because the unidirectional nonreciprocal coupling causes one of the modes to be unstable and only one mode survives. For instance, \(\Delta \phi = -\pi/2\) corresponds to a unidirectional forward coupling, meaning that the energy irreversibly flows from \(a_m\) to \(a_{m+1}\) and, thus, only \(a_{m+1}\) survives as the eigenmode (\(\beta_m = 1\)). It is worth noting that the results...
revealed that forward transmission is allowed (Fig. 5a, c), with adiabatic elimination (curves) agree well with those obtained by diagonalizing the effective Hamiltonian (2) and restricted time-reversal transformation. The residual imperfect adiabatic elimination of the connecting modes, whereas backward transmission is forbidden. The residual backward transmission, which originates from the imperfect adiabatic elimination of the connecting modes, is only $10^{-4}$ for $M = 2$ and $10^{-8}$ for $M = 3$ (Fig. 5b, d). The results calculated from the effective Hamiltonian (2) with adiabatic elimination (curves) agree well with those obtained from the original Hamiltonian (1). The insets illustrate the energy flow direction. The parameters are the same as in Fig. 4, with $\Delta \phi = -\pi/2$.

A unidirectional nonreciprocal coupling directly gives rise to unidirectional energy transmission between the main resonance modes. In Fig. 5, we plot the typical results of unidirectional energy transmission for two ($M = 2$) and three ($M = 3$) main modes in the case of a forward unidirectional coupling ($|h_+| \neq 0$ and $|h_-| = 0$). It is revealed that forward transmission is allowed (Fig. 5a, c), whereas backward transmission is forbidden. The residual backward transmission, which originates from the imperfect adiabatic elimination of the connecting modes, is only $10^{-4}$ for $M = 2$ and $10^{-8}$ for $M = 3$ (Fig. 5b, d). The results calculated from the effective Hamiltonian (2) with adiabatic elimination (curves) agree well with those obtained from the original Hamiltonian (1) (dots).

Discussion

Breaking Lorentz reciprocity is a very challenging task due to the time-reversal symmetry and linear nature of Maxwell’s equations, and at present, only three routes towards realizing nonreciprocity have been discovered, using a magnetic field, spatiotemporal modulation, or nonlinearity. Although loss breaks the time-reversal symmetry according to the traditional definition of time reversal, it is commonly believed that in classical electrodynamics, only restricted time reversal is valid. In this framework, lossy materials remain lossy under time-reversal transformation and restricted time-reversal symmetry still holds, which implies reciprocity$^{2,3}$. Here we demonstrate that nonreciprocity can be achieved by making use of loss combined with multichannel interference. Due to the interference with different loss angles, the field ratios do not remain conserved under time-reversal transformation.

Although PT-symmetric systems also make use of losses$^{8,10,48}$, it is clear that the nonreciprocity reported in previous PT-symmetric cavity schemes originates from the nonlinear gain saturation effect, i.e., the basic principle used to generate nonreciprocity is nonlinearity$^{9,10}$. In our scheme, we focus on a pure lossy and linear system, in which neither gain nor nonlinearity is required and the fundamental aspect giving rise to nonlinearity is loss. It is noteworthy that loss is ubiquitous, whereas gain and nonlinearity are not common in optical systems. In addition, we clarify that the above results represent nonreciprocity but not simply asymmetrical power transmission, as the input and output channels both contain a single mode and the scattering matrix is asymmetric (see the Supplementary Information for details), indicating the breaking of Lorentz reciprocity. This can also be verified when we consider the implementation of our scheme by means of an experimentally feasible setup in which single-mode standing-wave photonic crystal cavities are connected by waveguides$^{49,50}$ (see also the Supplementary Information for more details). The forward (backward) energy transmission coefficient can be defined as $T_- \equiv \left| \frac{a_{m+1}^{\text{out}}}{a_m^{\text{in}}} \right|^2$, which is equal to the modular square of the off-diagonal element of the scattering matrix and proportional to the modular square of the forward (backward) coupling coefficient $h_+ (h_-)$. Once the forward and backward coupling strengths are tuned to be unequal, the asymmetric scattering matrix leads to asymmetric forward and backward transmission coefficients, yielding nonreciprocity. Moreover, the unidirectional forward (backward) energy transmission coefficient can be maximized to 68.6% by optimizing the system parameters (the detailed derivations can be found in the Supplementary Information), corresponding to a 1.6 dB insertion loss.

In summary, we present the principle of loss-induced nonreciprocity, which is completely different from the existing principles relying on a magnetic field, spatiotemporal modulation, or nonlinearity. We design a coupled-mode model with a series of resonance modes interacting with each other via lossy connecting modes. A lossy mode possesses a phase lag induced by energy loss, which does not depend on the energy transmission direction. The interference between different coupling channels with different loss phases results in different coupling strengths for the forward and backward directions, yielding nonreciprocity. This property can be exactly tuned by matching the coherent coupling phases and loss phases,
which depend on the ratio between the detuning and energy decay rate of the resonance modes. Our model is universal and can be applied to a variety of systems that can be described by resonance modes, such as optical cavities and waveguides, mechanical resonators, and superconducting circuits. Our work provides new opportunities for designing nonreciprocal optical devices and exploring topological properties such as the non-Hermitian skin effect without requiring a magnetic field, spatiotemporal modulation, nonlinearity, or other stringent conditions, and it may also inspire the further exploration of methods of turning harmful effects into resources.

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Author contributions

Y.L. conceived the idea. Y.L. and X.H. carried out the calculations and wrote the paper. C.L. provided advice and helpful discussions. All authors contributed to the discussion of the results. Y.L. supervised the project.

Conflict of interest

The authors declare that they have no conflict of interest.

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