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**Chiral Selection of Purely Localized and Delocalized States in Non-Hermitian
Photonic Waveguide Arrays**

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Abstract

Advances in topological and non-Hermitian photonics, together with localization physics in aperiodic lattices, have expanded on-chip light-transport control, including the concentration and distribution of optical signals for integrated information flow. By hosting coexisting localized and delocalized states, systems with mobility edges provide a natural state basis for spatially confined and distributed outputs relevant to optical funneling and routing. However, this coexistence alone does not enable deterministic selection between the two output forms. Here, we design a non-Hermitian quasiperiodic mosaic waveguide array that integrates localized–delocalized coexistence with exceptional-point encircling dynamics, enabling deterministic chiral selection of a purely localized or delocalized output state. The engineered dissipative coupling generates the exceptional points and shapes the relevant eigenvalue branches with distinct eigenstate spatial profiles. A properly designed loop then makes the encircling direction select the localized or delocalized state on the lowest-loss energy surface. For experimental validation, we fabricate a waveguide-based chiral beam splitter-combiner in glass using femtosecond laser direct writing. Our results reveal an approach that combines non-Hermitian encircling dynamics with localization physics, enabling precise control of on-chip light transport.

Introduction

Integrated photonic waveguide systems provide a scalable on-chip platform for communication and information processing, where on-demand light manipulation and routing are central goals. Topological and non-Hermitian photonics have enriched wave-transport control through topological protection and exceptional-point dynamics^{1,2}, while localization physics offers additional control by

reshaping eigenstate spatial distributions and determining whether light spreads or remains confined during propagation^{3,4}. Despite these advances, integrated photonic systems still lack a mechanism for chiral selection between purely localized and delocalized states. Such functionality is important for on-chip information flow like the funneling and routing because it allows the propagation direction to select between the spatial concentration and distribution of optical signals.

Realizing the chiral localized-delocalized state selection (LDSS) requires a system that supports both spatially concentrated and distributed optical outputs. This motivates the coexistence of localized and delocalized states within a single platform. Systems with mobility edges (MEs), which separate localized from delocalized states in energy space⁵, provide a suitable platform for this purpose. While MEs are a hallmark of three-dimensional random disorder systems⁶, they can also emerge in certain one-dimensional quasiperiodic systems through tailored hopping processes⁷⁻¹⁰ or suitably modified quasiperiodic potentials¹¹⁻¹⁴. In such systems, the energy-dependent localization-delocalization transition can be driven by altering a parameter beyond a critical point, thereby enabling controllable transport properties¹⁵.

However, the coexistence of localized and delocalized states alone is insufficient for realizing chiral LDSS. A mechanism that can steer state evolution toward the desired output is also needed. Non-Hermitian physics can introduce additional dynamical and spectral degrees of freedom for state manipulation¹⁶⁻¹⁹, and its interplay with quasiperiodicity has recently attracted growing attention²⁰⁻²⁷. In particular, dissipation has been shown to induce localization-delocalization transition in systems with MEs^{28,29}. By tuning a specific dissipative phase, the system can be driven into a steady state dominated by either localized or delocalized components. Yet these schemes do not provide directional selection, and their outputs are generally mixed steady states rather than unambiguous

65 purely localized or purely delocalized states. Together with the practical difficulty of implementing
66 such phase-engineered Lindblad dissipation in integrated waveguide arrays^{30–32}, this motivates the
67 search for a different non-Hermitian approach.

68 Dynamically encircling exceptional points (DEEPs) provides a distinct chiral state-selection
69 method, where the final state is determined by the encircling direction and is independent of the
70 initial state^{33–38}. This stems from the nonadiabatic transitions^{39,40} due to the breakdown of the
71 adiabatic theorem during dynamic evolution in non-Hermitian systems, driving the final state to
72 reside on the lowest-loss energy surface. Toward chiral LDSS, the key task is therefore to construct
73 a photonic lattice that combines localized–delocalized coexistence with DEEP dynamics.

74 In this paper, we realize the chiral LDSS in a class of non-Hermitian quasiperiodic mosaic
75 systems processing DEEPs. The introduced dissipative coupling⁴¹ induces multiple pairs of EPs in
76 the energy spectrum and gives rise to spectral branches whose corresponding eigenstates exhibit
77 different spatial localization. The encircling trajectory is designed in a two-dimensional parameter
78 space spanned by the coupling strength and relative dissipation. The two segments of the loop
79 extending from the starting point in opposite directions traverse regions associated with localized
80 and delocalized states, respectively, on the lowest-loss energy surface. Therefore, the propagation
81 direction then deterministically selects the final state to be either purely localized or purely
82 delocalized, irrespective of the initial excitation. It is verified by a chiral beam splitter-combiner
83 composed of a waveguide array fabricated inside glass using the femtosecond laser direct writing^{42–}
84 ⁴⁴. Our results provide a new approach to controllable light transport in integrated waveguide
85 platforms and open further avenues for developing advanced photonic devices.

Non-Hermitian quasiperiodic mosaic model:

The one-dimensional quasiperiodic mosaic lattice is a class of exactly solvable models that host MEs in their spectra¹³. Photonic mobility-edge phases based on this model have been experimentally observed in femtosecond-laser-written waveguide lattices¹⁵. Yet the Hermitian version lacks the necessary conditions to realize DEEPs. We therefore extend this model to the non-Hermitian regime by the addition of the non-Hermitian quasiperiodic onsite potential Δ_n to the Hamiltonian:

$$H = C \sum_n^N (a_n^\dagger a_{n+1} + H.c.) + \sum_n^N \Delta_n a_n^\dagger a_n,$$

where $a_n(a_n^\dagger)$ is the annihilation (creation) operator at site n , C denotes the nearest-neighbor coupling strength, and Δ_n are given by:

$$\Delta_n = \begin{cases} 2\Delta \cos(2\pi bn + \phi) + i\gamma_q, & n = mk, \\ i\gamma_z, & \text{otherwise.} \end{cases}$$

with $b = (\sqrt{5} - 1)/2$. Here, k is an integer that controls the spatial positions at which the quasiperiodic potential is applied. The parameters Δ , ϕ , γ_q , γ_z denote the quasiperiodic potential amplitude, the phase offset, the intrinsic loss on quasiperiodic potential sites and on relative zero potential sites, respectively. The localization is characterized by the inverse participation ratio (IPR):

$$IPR_i = \left(\sum_n^N |\psi_i^n|^4 \right) / \left(\sum_n^N |\psi_i^n|^2 \right)^2,$$

where the superscript i denotes the i th eigenstate, and n labels the lattice site. The IPR ranges from $1/N$ for fully delocalized states to 1 for fully localized states.

To ensure compatibility with the waveguide-array implementation, we choose parameters within experimentally feasible ranges, and fix $\Delta = 0.2 \text{ mm}^{-1}$, $\phi = 0.2\pi$, $\gamma_q = \gamma_z = 0.2 \text{ mm}^{-1}$ and $k = 2$, yielding the model shown in Fig. 1a. For conservative coupling $C = J (J \in \mathbb{R})$, the spectrum remains essentially the same as in the Hermitian case, except for an overall shift in its imaginary part, as shown in Fig. 1b. The mobility edges are located at $E_{ME} = \pm J^2/\Delta$, separating spatially

localized eigenmodes from delocalized ones¹⁵. When the coupling is instead dissipative $C = iJ$ ($J \in \mathbb{R}$), multiple pairs of exceptional points emerge in the spectrum, as shown in Fig. 1c. For a pair of energy branches that coalesce at an EP located at J_{EP} , the real parts of the eigenvalues bifurcate for $J < J_{EP}$, while the corresponding eigenstates on the two branches exhibit different *IPR* under identical system parameters. This differs from conventional two-state PT- or anti-PT-symmetric systems, where eigenstates on the two branches typically share the same *IPR* under identical system parameters. We exploit the spectral bifurcation and the different degrees of spatial localization of the corresponding eigenstates to realize a chiral LDSS via DEEPs.

The designed position-dependent evolution trajectories for $J(z)$ and $\gamma(z)$ form a closed loop winding around EPs in the $J - \gamma$ parameter space, where $\gamma = \gamma_q - \gamma_z$. For any given (N, k) , there exists a unique lowest-loss energy surface. On this surface, the calculated *IPR* of the eigenstates reveals asymmetry: fully localized states exist only for $\gamma \leq 0$, in sharp contrast to the symmetric *IPR* distribution in conventional two-state PT- or anti-PT-symmetric systems (see Fig. 1d). Since the final state in the DEEP process is ultimately selected to lie on the lowest-loss energy surface³⁶, an appropriately designed closed loop is required for chiral selection of a localized or delocalized output state. We configure the encircling trajectory in the $J - \gamma$ parameter space so that its two segments extending from the starting point in opposite directions pass through high- and low-*IPR* regions of the lowest-loss energy surface, corresponding to localized and delocalized states, respectively. Only under this condition do the two eigenstates associated with the starting-point parameters on the lowest-loss energy surface consist of one localized and one delocalized state, both serving as direction-selectable outputs during the dynamical evolution.

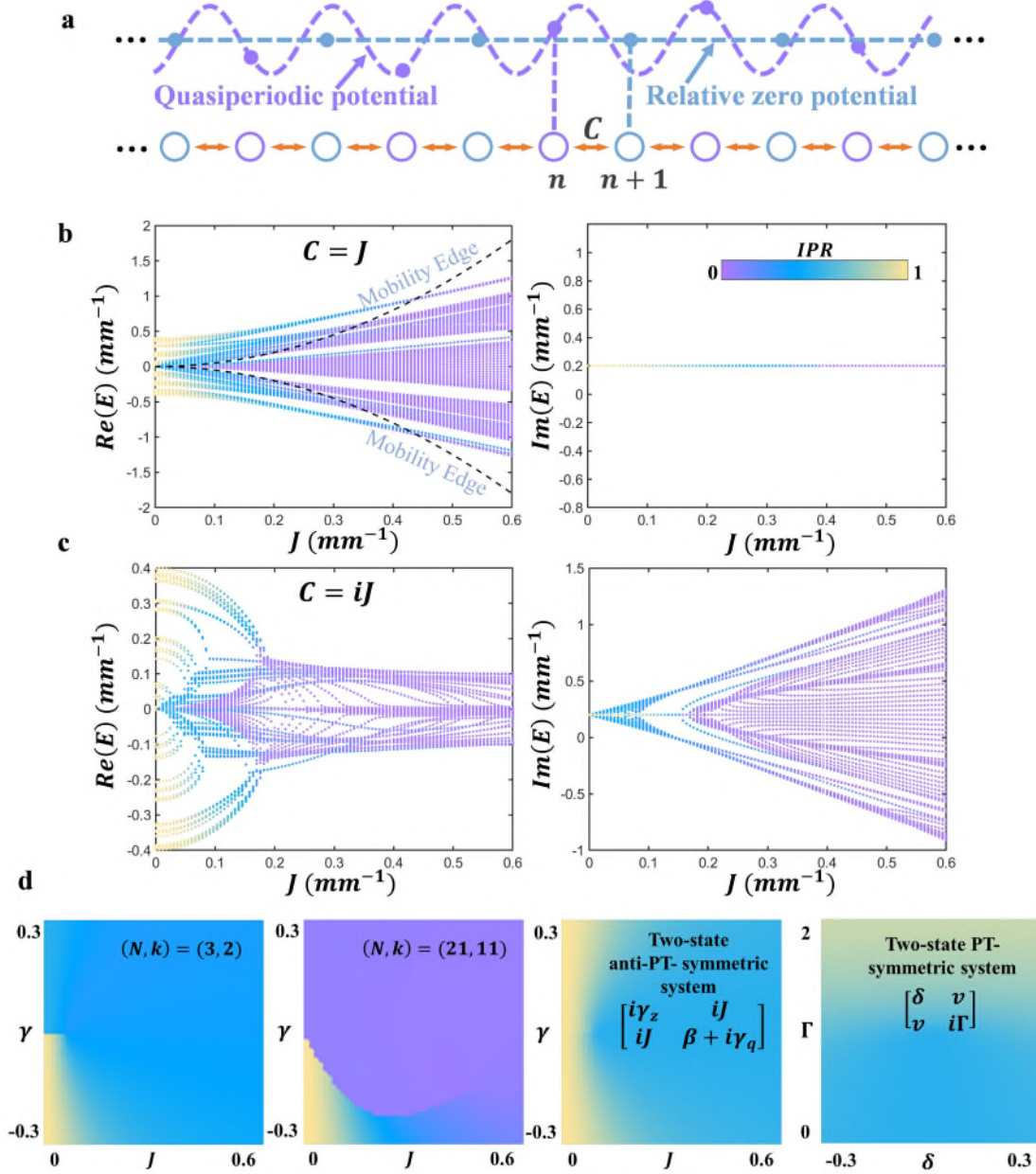


Fig. 1. **a** The non-Hermitian quasiperiodic mosaic model with $k = 2$. The purple and blue hollow circles denote sites with quasiperiodic and relative-zero potentials, respectively; the filled dots on the dashed curves mark their corresponding potential values. C denotes the nearest-neighbor coupling strength. **b**, **c** The real and imaginary parts of the eigenvalue spectrum for conservative coupling $C = J$ (**b**) and dissipative coupling $C = iJ$ (**c**). The color denotes the inverse participation ratio (IPR) of the corresponding eigenstate. The dashed lines in **b** represent the theoretically predicted mobility edges. The parameters we set are $\Delta = 0.2 \text{ mm}^{-1}$, $\phi = 0.2\pi$, $\gamma_q = \gamma_z = 0.2 \text{ mm}^{-1}$, $k = 2$ and $N = 89$. **d** IPR

of eigenstates on the lowest-loss energy surface for different (N, k) , compared with that of a two-state PT- and anti-PT-symmetric system.

Chiral LDSS through DEEPs:

We first consider the case $k = 2$. For $N = 3$, the real and imaginary parts of the eigenvalues in the $J - \gamma$ parameter space are shown in Fig. 2a, where the lowest-loss energy surface is highlighted in sky blue. The encircling trajectory starts at $J = 0.01 \text{ mm}^{-1}$ and $\gamma = 0$, a point characterized by three eigenstates ψ_1 , ψ_2 and ψ_3 whose spatial distributions are shown in Fig. 2b. The eigenstates ψ_1 and ψ_2 reside on two energy surfaces that are connected at an EP located at $J = \Delta \cos(4\pi b) / \sqrt{2}$, $\gamma = 0$.

We find that the DEEP process enables direction-dependent selection between the two target eigenstates ψ_1 and ψ_2 . For clockwise (CW) encircling (Fig. 2c), when the input state is ψ_2 , the highest occupied state remains on the lowest-loss energy surface throughout the evolution (illustrated by the blue trajectory), yielding ψ_1 as the final output state. When the input state is ψ_1 , the highest occupied state initially evolves on the highest-loss energy surface and subsequently transitions to the lowest-loss surface through the nonadiabatic transition process; the final output state is again ψ_1 . These results demonstrate that, for CW encircling, the output state is always ψ_1 . For counterclockwise (CCW) encircling, the output state is always ψ_2 (Fig. 2d). The dynamics of the normalized intensity evolution at each site for both CW and CCW encircling processes are shown in Fig. 2e. The calculated IPR of ψ_1 and ψ_2 clearly distinguish the two states: $IPR_1 = 0.991$ confirms that ψ_1 is a localized state, whereas $IPR_2 = 0.495$ indicates that ψ_2 is delocalized. This minimal sparse-potential lattice model preserves the essential localized–

delocalized coexistence, and the direction-dependent state selectivity of DEEPs further enables
chiral LDSS.

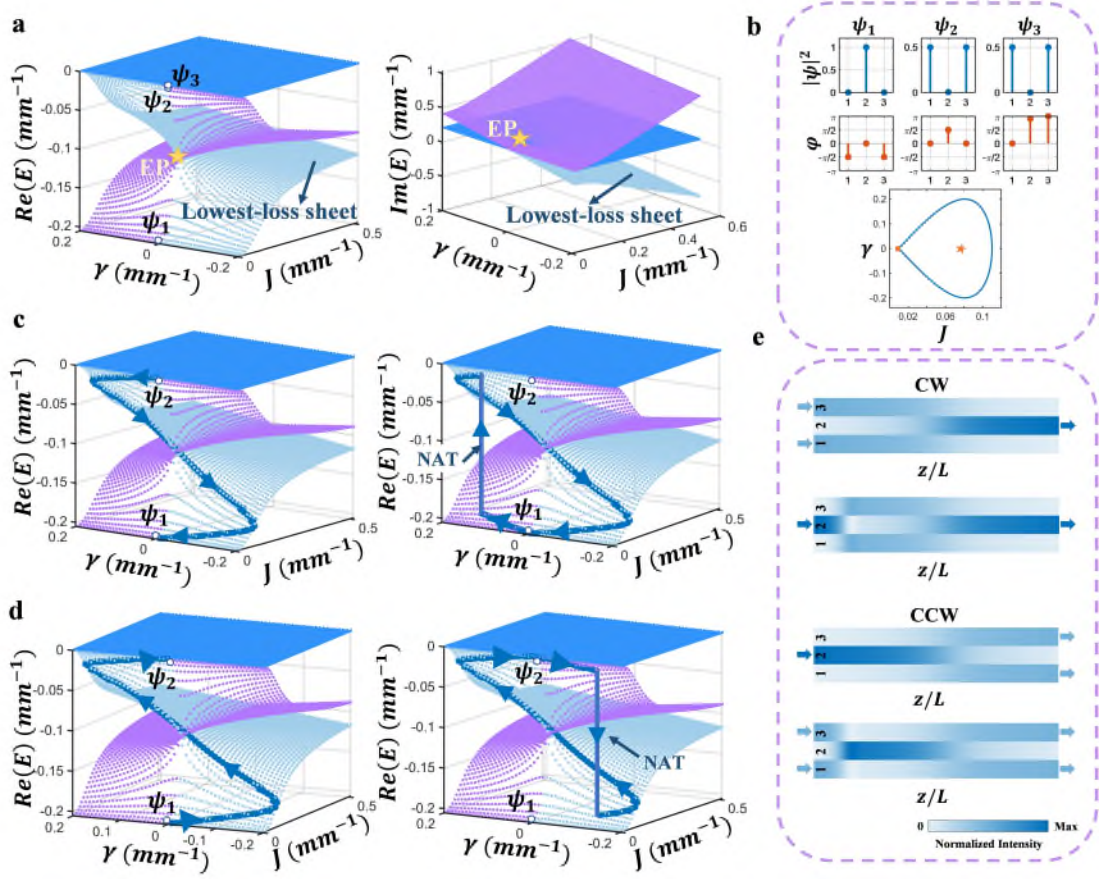


Fig. 2. Simulated results of the chiral LDSS for $k = 2$ and $N = 3$. **a** The real and imaginary parts of energy surfaces in the $J - \gamma$ parameter space. **b** The encircling trajectory and the distribution of eigenstates at the starting point. **c, d** Mode conversion results for the CW (**c**) and CCW (**d**) encircling direction with input modes ψ_1 and ψ_2 . **e** The dynamics of the normalized intensity evolution at each site for CW (top) and CCW (bottom) encircling processes.

As N increases, chiral LDSS between states with different IPR remains achievable. However, while one state in each pair is fully localized ($IPR = 1$), the other does not correspond to a fully delocalized mode ($IPR = 1/N$). Instead, its wavefunction extends over only a limited number of

waveguides (see Supplementary information, Sec. I). Fortunately, we find, by adjusting k accordingly as N increases, it is still possible to realize the direction-dependent selection between a fully localized state and a state approaching full delocalization. Guided by the IPR scaling and the spatial profiles of the eigenstates, we choose $k = (N + 1)/2$ for odd N , which yields chiral LDSS between a fully localized state ψ_L and a delocalized state ψ_D whose wavefunction is well described by a double-Gaussian profile.

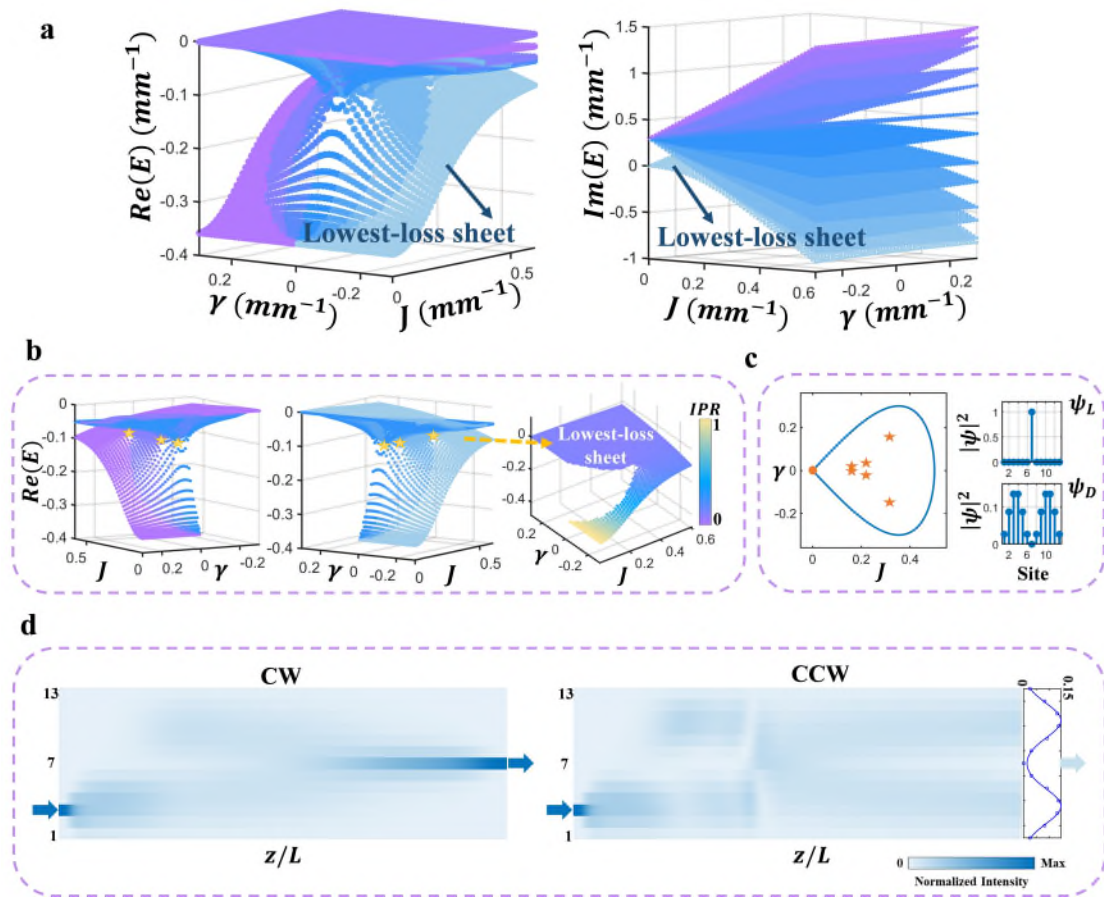


Fig. 3. Simulated results of the chiral LDSS for $k = 7$ and $N = 13$. **a** The real and imaginary parts of energy surfaces in the $J - \gamma$ parameter space. **b** Left and middle panels: the six energy surfaces with nonzero real parts hosting six EPs (yellow stars). The remaining seven energy surfaces, which are purely imaginary and exhibit no EPs, are omitted for clarity. Right panel: the calculated IPR of the eigenstates on the lowest-loss energy surface. **c** The encircling trajectory and the spatial distribution of ψ_L and ψ_D .

d Dynamics of the normalized intensity evolution for CW (left) and CCW (right) encircling processes, with initial excitation at site 3.

We present a detailed analysis for the parameters $(N, k) = (13, 7)$. The real and imaginary parts of the energy surfaces in the $J - \gamma$ parameter space are shown in Fig. 3a. Figure 3b displays the six EPs (marked by yellow stars) hosted in the system, together with the calculated IPR of the eigenstates on the lowest-loss energy surface. Accordingly, an encircling trajectory is designed following the aforementioned principles, as shown in Fig. 3c. The trajectory starts at $J = 0.0001 \text{ mm}^{-1}$ and $\gamma = 0$, with the spatial distributions of ψ_L and ψ_D shown in the insets. ψ_L is fully localized at the central site ($IPR_L = 1$), whereas ψ_D fits a double-Gaussian distribution ($IPR_D = 0.107$). For the CW encircling direction, the output state is always ψ_L , independent of the initial state. For the CCW encircling direction, the system always evolves into ψ_D , irrespective of the initial excitation. The dynamics of the normalized intensity evolution at each site for CW and CCW encircling processes, with site 3 initially excited, are shown in Fig. 3d. Additional results for other excitation conditions and different (N, k) values are provided in Supplementary information, Sec. II.

Experimental realization:

The chiral LDSS realized here enables the direction-controlled switching between signal collection and distribution, providing a physical basis for information flow like optical funneling and routing. As an experimental realization, we further demonstrate a chiral beam splitter–combiner in a photonic waveguide system fabricated by femtosecond laser direct writing. The corresponding

effective Hamiltonian is given by

$$H_1(z) = \begin{bmatrix} i\gamma_z & iJ(z) & 0 \\ iJ(z) & \beta + i\gamma_q(z) & iJ(z) \\ 0 & iJ(z) & i\gamma_z \end{bmatrix},$$

with $\beta = 2\Delta\cos(4\pi b + \phi)$. Since dissipative coupling cannot be directly implemented in photonic

waveguide platforms, we instead utilize the following five-state Hamiltonian:

$$H_2(z) = \begin{bmatrix} \beta_1 & \kappa_1(z) & 0 & 0 & 0 \\ \kappa_1(z) & \beta_2 + i\gamma & \kappa_2(z) & 0 & 0 \\ 0 & \kappa_2(z) & \beta_3 & \kappa_2(z) & 0 \\ 0 & 0 & \kappa_2(z) & \beta_2 + i\gamma & \kappa_1(z) \\ 0 & 0 & 0 & \kappa_1(z) & \beta_1 \end{bmatrix},$$

where β , κ , and γ denote the propagation constant, coupling strength, and loss of waveguides,

respectively. By adiabatically eliminating the auxiliary lossy modes, the five-state system is shown

to be equivalent to the three-state model H_1 , but the high loss required for adiabatic elimination

causes non-negligible evolution path-dependent losses, which reduce the transmission efficiency to

24.7% (see Supplementary information, Sec. III).

To achieve relatively high transmission efficiency, we further optimize $H_2(z)$ to reduce path-

dependent losses. We set the detuning between all sites to zero, leading to the following Hamiltonian:

$$H_3(z) = \begin{bmatrix} 0 & \kappa_1(z) & 0 & 0 & 0 \\ \kappa_1(z) & i\gamma & \kappa_2(z) & 0 & 0 \\ 0 & \kappa_2(z) & 0 & \kappa_2(z) & 0 \\ 0 & 0 & \kappa_2(z) & i\gamma & \kappa_1(z) \\ 0 & 0 & 0 & \kappa_1(z) & 0 \end{bmatrix}.$$

The eigenvalues of the system are $E_1 = 0$, $E_{2,3} = (i\gamma \mp \sqrt{4\kappa_1^2 - \gamma^2})/2$, $E_{4,5} = (i\gamma \mp$

$\sqrt{4\kappa_1^2 + 8\kappa_2^2 - \gamma^2})/2$, whose real and imaginary parts in the $\kappa_1 - \kappa_2$ parameter space are shown

in Fig. 4a. The system exhibits two exceptional lines (ELs) located at $\kappa_1 = \gamma/2$, $\kappa_1^2 + 2\kappa_2^2 = \gamma^2/4$,

respectively. Notably, the imaginary part of the lowest-loss energy surface of the system is zero,

thereby significantly reducing path-dependent losses. Under such conditions, high-efficiency

direction-dependent state selection can be realized based on dynamically encircling ELs by properly

designing the encircling trajectory. Here, we set $\gamma = 0.4 \text{ mm}^{-1}$ and vary κ_1 and κ_2 along the trajectory shown in Fig. 4b. The simulated state evolution yields a transmission efficiency of 78.1%. This value is limited by the feasible coupling strength κ and the total evolution length in the experiment. It can approach unity when either parameter increases sufficiently (see Supplementary information, Sec. III).

We map the encircling trajectory onto the coupling distance between waveguides in the EL encircling region, which forms part of the coupled waveguide system shown in Fig. 4c and is sandwiched between two auxiliary regions designed for optical fan-in and fan-out. The total device length is 80 mm. The experimentally measured output results for the CW and CCW encircling processes are shown in Fig. 4d. For the CW process, the five-waveguide system consistently outputs through WG3 regardless of input state, thereby functioning as an optical combiner. Conversely, under the CCW process, the output is nearly equally distributed between WG1 and WG5, regardless of the input, performing beam-splitting. The experimentally measured transmission efficiency, defined as the maximum transmittance among all conversion processes, is 59.4%. The discrepancy between simulation and experiment is mainly attributed to fabrication imperfections in the discrete scatterer loss implementation, which lead to nonuniform loss and thus affect the transmission efficiency (see Supplementary information, Sec. VI).

Numerical simulations further demonstrate that the EL encircling process is robust against local random perturbations in κ and γ , provided that the encircling trajectory does not undergo significant deformation (see Supplementary information, Sec. VII). The system therefore constitutes a robust chiral splitter-combiner capable of chiral information collection and distribution. In Secs. IV and V of the Supplementary information, we present beam propagation method simulations that

directly visualize the optical field evolution in a realistic five-waveguide system, along with an implementation of the $(N, k) = (5, 3)$ case using a nine-waveguide system, further validating the practical feasibility of our design.

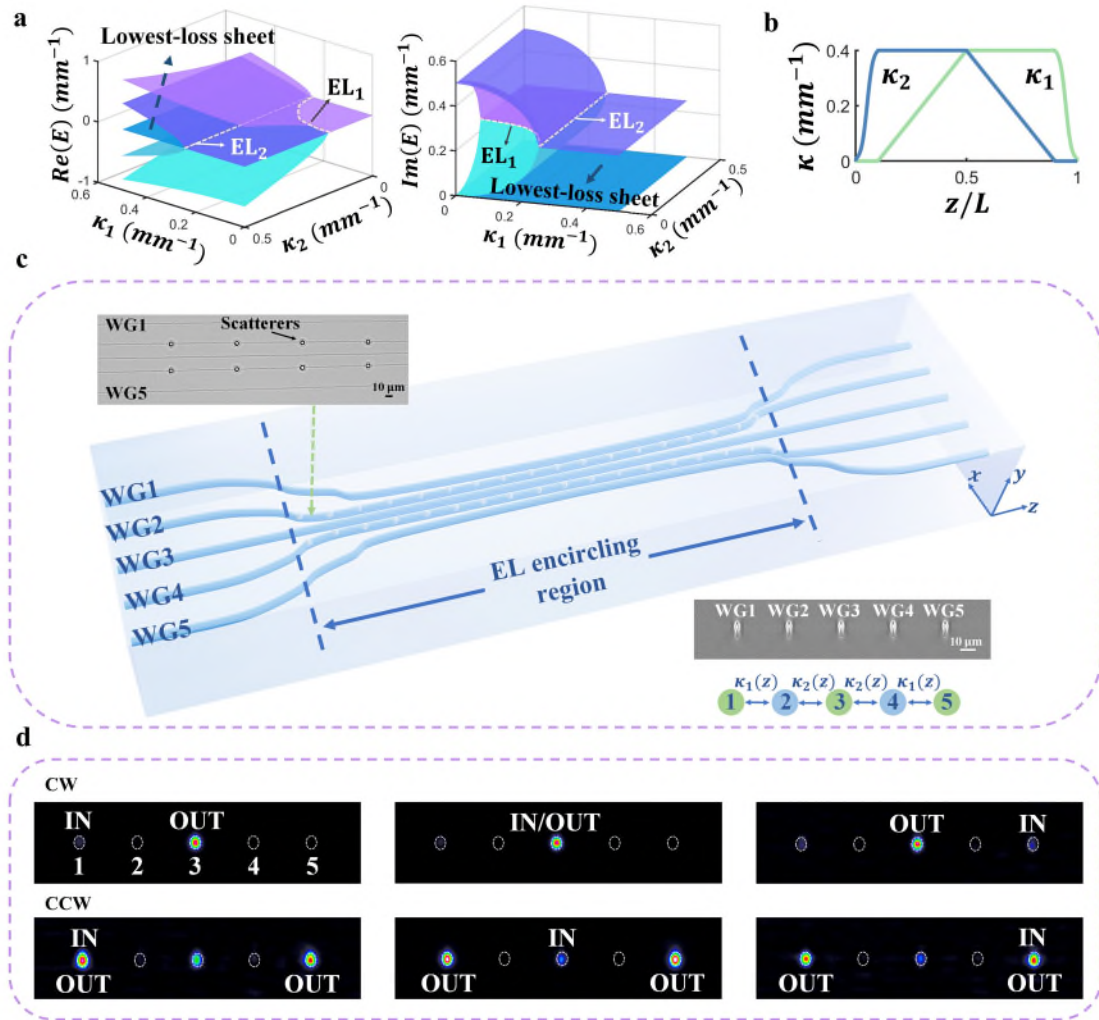


Fig. 4. Experimental realization of the chiral beam splitter-combiner. **a** The real and imaginary parts of the energy surfaces in the $\kappa_1 - \kappa_2$ parameter space at $\gamma = 0.4 \text{ mm}^{-1}$. **b** Evolution trajectories of κ_1 and κ_2 along the CW encircling direction. **c** Schematic diagram of the chiral beam splitter-combiner structure. Uniform loss is introduced in WG2 and WG4 by creating scatterers. Inset: Microscopic photographs of the fabrication sample. **d** Experimental results of the chiral mode conversion at 785 nm for the CW (top) and CCW (bottom) encircling direction.

258

259 **Discussion:**

260 In summary, we demonstrate DEEP-enabled chiral selection between purely localized and purely
261 delocalized transport outcomes. The propagation direction dictates whether light in an on-chip
262 photonic system is funneled toward a single output port or redistributed into a delocalized state. This
263 enables chiral signal collection and distribution on chip, providing a physical basis for optical
264 funneling and routing. The fabricated five-waveguide non-Hermitian array functions as a chiral
265 beam splitter–combiner, providing direct experimental validation of the proposed mechanism.
266 Moreover, the mechanism suggests a straightforward extension to active platforms. With suitable
267 gain media, the pump-profile modulation could be used to switch between laser supermodes with
268 distinct IPR ⁴⁵ to realize robust and fast reconfigurable on-chip lasers. Notably, the underlying
269 mechanism is not restricted to photonic waveguide platforms and can, in principle, be generalized
270 to other physical settings, including atomic and electronic systems. Our results provide a novel
271 approach to manipulating transport and localization properties, opening up new pathways for
272 enriched mode control and novel information-processing functionalities.

273

274 **Materials and methods:**

275 The photonic waveguide systems are fabricated in borosilicate glass (Eagle XG, Corning) by
276 focusing pulses generated by a regeneratively amplified Yb: KGW femtosecond laser system
277 (Pharos20, Light Conversion) at a wavelength of 1030 nm with a duration of 217 fs, and a 1 MHz
278 repetition rate. The microscope objective (20 ×, NA = 0.5) focuses the pulse with a power of 370
279 nJ at a depth of 170 μm below the glass surface to write the waveguides of the five-waveguide

structure (Fig. 4c). We move samples at a speed of 25 mm/s by using a computer-controlled high-precision three-axis air-bearing stage (FG1000, Aerotech). The fabricated waveguide systems support single-mode propagation at 785 nm. To introduce additional losses into WG2 and WG4 in the five-waveguide structure, we focused laser pulses at uniformly spaced points within the waveguides, with each point separated by 100 μm . The pulse energy to create scatterers is 430 nJ, and the exposure time is 0.5 s. Each scatterer was measured to induce a loss of approximately 0.3 dB.

Data availability:

The data that support the findings of this study are available from the corresponding author upon reasonable request.

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R.Q.W. and Y.L. conceived of the idea. R.Q.W. performed the theoretical analysis and conducted the numerical simulations. R.Q.W. and Y.L.T. performed the experiments. The manuscript was written by R.Q.W. The project was supervised by Q.H.G. and Y.L. All authors discussed the results and contributed to the manuscript.

Conflict of interest:

The authors declare no competing interests.

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