Chiral control of quantum states in non-Hermitian spin-orbit-coupled fermions

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Spin–orbit coupling is an essential mechanism underlying quantum phenomena such as the spin Hall effect and topological insulators. It has been widely studied in well-isolated Hermitian systems, but much less is known about the role of dissipation in spin–orbit-coupled systems. Here we implement dissipative spin–orbit–coupled bands with ultracold fermions, and observe parity-time symmetry breaking as a result of the competition between the spin–orbit coupling and dissipation. Tunable dissipation, introduced by state-selective atom loss, enables us to tune the energy gap and close it at the critical dissipation value, the so-called exceptional point. In the vicinity of the critical point, the state evolution exhibits a chiral response, which enables us to tune the spin–orbit coupling and dissipation dynamically, revealing topologically robust chiral spin transfer when the quantum state encircles the exceptional point. This demonstrates that we can explore non-Hermitian topological states with spin–orbit coupling.

An open quantum system that does not conserve energy is effectively described by a non-Hermitian Hamiltonian and exhibits various counterintuitive phenomena that cannot exist when the system is Hermitian. One such example is the fundamental understanding of non-Hermitian topological matter that may require subtle classification in contrast to the Hermitian case, such as iconic topological insulators. Although extensive theoretical1–4 and experimental5–9 works have been carried out, how the non-Hermitian topological state can be classified remains elusive. In particular, spin–orbit coupling (SOC), a key mechanism driving the non-trivial band topology, has not been investigated in non-Hermitian quantum systems. Recently, considerable efforts have been made in ultracold atoms to explore synthetic SOCs10–13 and the associated topological bands14–18, and therefore ultracold atoms offer the unprecedented possibility of studying the non-Hermitian SOC mechanism, a critical step towards realizing non-Hermitian topological phases.

In this Letter we make an important step in this direction by realizing non-Hermitian spin–orbit–coupled fermions and observing a parity-time (PT) symmetry-breaking transition as a result of the competition between SOC and dissipation. We implement synthetic SOC for ultracold fermions10 together with non-Hermiticity, tunable in time. This controllability allows us to investigate how the spin–orbit-coupled energy band changes with dissipation and to explore the PT symmetry-breaking transition across the exceptional point (EP) and its topological nature. Exploring the parameter regime from SOC-dominated to loss-dominated behaviour, we identify an EP in a fully quantum regime. Finally, we experimentally probe the chiral property of the EP by encircling it in a parameter space and showing chiral quantum state transfer due to the breakdown of adiabaticity. Our work sets the stage for the experimental study of many-body states in the complex energy bands across the PT symmetry-breaking transition. Recently, the feasibility of realizing unprecedented phenomena in non-Hermitian atomic systems was noted, including enhanced fermionic superfluids19, the generalized Floquet time crystal with dissipation20 and higher-order topological phases21. This work complements the non-Hermitian phenomena observed in classical systems, including topological energy transfer22–24, enhanced sensing25, PT symmetric lasing26 and novel topological entities such as exceptional rings27 and non-Hermitian topological edge states28. Nevertheless, the role that non-Hermiticity plays in the genuine quantum regime ranging from few-body to many-body systems remains largely unexplored. Although recent works have explored non-Hermitian systems within a quantum framework29 and have demonstrated the PT symmetry-breaking transition in non-Hermitian quantum systems, such as photons29,30, superconducting qubits31, single electronic spins32, exciton–polaritons33 or atoms19, the dynamic evolution of the quantum state near the EP has been limited to the single-particle quantum system.

Here we demonstrate the chiral control of the quantum state with ultracold fermions, showcasing many-particle quantum systems. Our experiment begins by loading ultracold fermions of 173Yb into the engineered energy band in which two hyperfine states (labelled $|\uparrow\rangle$ and $|\downarrow\rangle$) are coupled by a pair of Raman transition beams, resulting in synthetic SOC. In a typical energy dispersion, the energy gap between two dressed bands is opened by Raman two-photon detuning. Here we define the natural units of momentum and energy as $\hbar k = \sin(\theta/2) \frac{2\pi}{4556}$ and $E_i = \frac{\hbar^2 k_i^2}{2m} = h \times 1.41$ kHz, where $\lambda_{156} = 556$ nm. The real momentum $k_i$ is then related to the quasimomentum along the $\hat{x}$ direction as $k_i = q_i \mp k_{z}$ for spin–up and spin–down, respectively. In our experiment, atom loss is induced by the single near-resonant beam (Fig. 1b), resulting in the fixed ratio of $\gamma_{1}/\gamma_{2} = 13$.

The quantum dynamics of two dressed energy bands $\{|+\rangle, |-\rangle\}$, corresponding to the eigenvalues $\lambda_{\uparrow}$ of $\mathcal{H}$, are governed by the eigenvalue difference $\Delta \lambda = (\lambda_{\uparrow} - \lambda_{\downarrow})$. In the $q_i–\gamma_{1}$ parameter space, the
\[ \delta = \text{symmetric phase, the initial quantum state oscillates} \]

\[ \delta \delta_3 \]

atoms in \( |\uparrow\rangle \) without the Raman transition and loss. Subsequently, \( |\uparrow\rangle \) atoms are transferred to the \( |\downarrow\rangle \) state when the Raman coupling is switched on (Fig. 1c). Here, brief pulses of the Raman and loss beams are applied for a variable time duration. Then a 6-ms time-of-flight expansion maps the momentum to real space, allowing direct momentum resolution of the energy band.

One of the features associated with dissipation is the closing of the energy gap at the EP\(^+\). Figure 2a shows typical spin–orbit-coupled energy bands at different dissipation strengths for \( \delta = 4E \), and their Rabi oscillations persisting in the Raman field with dissipation for a variable duration (Fig. 2b). In the absence of atom loss (\( \gamma_2 = 0 \)), the state undergoes Rabi oscillations associated with the energy gap at each quasimomentum. The Rabi coupling is resonant at \( q_x = 1.0k_r \), revealing the smallest energy gap. The Rabi oscillation becomes slower with increasing dissipation. By fitting the spin oscillation with a damped sinusoidal function, we determine the energy gap and damping rate at each quasimomentum as described in Fig. 2b.

A Fermi sea collectively undergoes Rabi oscillations, in contrast to the non-Hermitian classical or single-particle quantum\(^{\text{–}}\) system, revealing a non-uniform bandgap in time-dependent spin textures (Fig. 3a). From this quasimomentum-dependent Rabi oscillation, we extract the energy gap, \( \text{Re}(\Delta \lambda) \), between two dressed bands for different dissipation strengths, as shown in Fig. 3b, in good agreement with the theoretical expectation. The dissipation gives rise to gap closing at the resonant coupling at \( q_x = 1.0k_r \), while the other quasimomenta still reveal a finite energy gap.

Figure 3d,e shows a quantitative measurement of the bandgap (\( \text{Re}(\Delta \lambda) \)) and damping rate (\( \text{Im}(\Delta \lambda) \)) of the non-Hermitian system at \( q_x = 1.0k_r \), manifesting the \( PT \) symmetry-breaking transition. In the \( PT \) symmetric phase, the initial quantum state oscillates between two eigenstates of the non-Hermitian Hamiltonian (Fig. 3c). Above the EP, however, the strong dissipation completely closes the energy gap with finite \( \text{Im}(\Delta \lambda) \neq 0 \), giving rise to monotonous behaviour. Both the bandgap and damping rate show the \( PT \) symmetry-breaking transition near the EP at \( q_x = 9.75E \) (Fig. 3d,e). At a non-resonant momentum, \( q_x = 0.5k_r \), the energy gap saturates to a finite value. With large dissipation, the spin–orbit-coupled band structure becomes similar to free particle dispersion, which may be related to the quantum Zeno effect (Methods\(^{\text{–}}\)).

We now explore topological features near the EP by dynamically encircling it with fermions. Near the EP, the state evolution is direction-dependent, resulting in intriguing chiral behaviour when encircling the EP due to the breakdown of adiabaticity in non-Hermitian systems\(^{\text{–}}\). Although this topological chirality has been observed in classical systems\(^{\text{–}}\), such as photonics and acoustics, it remains largely unexplored in quantum systems\(^{\text{–}}\), especially in many-body quantum systems. Here we dynamically encircle the EP with fermions in different directions. This is enabled by the EP occurring at the quasimomentum \( q_x = (\frac{\pi}{4})k_r \), where SOC is resonant. This leads to control of the EP position along the \( q_x \) axis, effectively with respect to the Fermi sea in the dressed band, by tuning \( \delta \). In our experiment, we simultaneously tune the loss (\( \gamma_r \)) and two-photon detuning (\( \delta \)) and dynamically explore arbitrary paths in the complex band.

We begin with adiabatic loading of fermions into the lowest energy band at \( \delta = 3E \). Initially, a Fermi sea is adiabatically formed at \( q_x = -0.8k_r \), with the dominant spin component of \( |\downarrow\rangle \) (blue colour of energy band, Fig. 4), whereas the EP occurs at \( q_x = (\frac{\pi}{4})k_r = 0.75k_r \), in the \( q_x - \gamma_r \) parameter space. We now tune \( \delta \) to \( +3E \), to \( -6E \), which shifts the EP to \( q_x = -1.5k_r \). Subsequently, the loss is increased from zero to different \( \gamma_r \) parameter space. We now follow by consecutive control of the two-photon detuning and loss. This results in counterclockwise encircling along a closed loop within \( T = 10.1 \) ms in the \( q_x - \gamma_r \) parameter space, as depicted in Fig. 4a (left), where \( T \) is the total encircling time. By reversing the aforementioned process, we can encircle the EP in a clockwise manner.
The spin polarization of the quantum state after the encircling is then determined via an optical Stern–Gerlach pulse followed by ballistic expansion. When the atom loss $r$ is increased to $r^{\text{max}} = 10E_r$ or $8E_r$, we observe that the initial quantum state $|\downarrow\rangle$ is selectively switched to spin-up depending on the encircling direction (Fig. 4b). This observation reveals that the evolution of the quantum state switches the state to a different eigenstate near the EP depending on the encircling direction. The chiral spin transfer disappears for $r^{\text{max}} = 2E_r$ if the EP is located far from the loop (Fig. 4c).

To better understand the quantum dynamics during non-adiabatic evolution, we performed numerical calculations in a larger domain of accessible parameters in $r^{\text{max}}$ and encircling duration. The state conversion efficiencies $C$ for encircling in the two directions are numerically calculated while considering each momentum sector integrated over the Fermi sea, and are plotted as two colour maps in Fig. 4a. The experimental results in Fig. 4b,c (labelled 1–4) are indicated by circles at $T = 10.1\,\text{ms}$, in good agreement with $C > 0$ for case 1 and $C = 0$ for the other cases. In fact, a smaller encircling speed (that is, a larger $T$) increases the chance of non-adiabatic transitions as long as the trajectory runs near the EP and further lowers the transition threshold of $r^{\text{max}}$ (ref. 23). The chiral behaviour is also observed for $T = 5\,\text{ms}$. On the other hand, we observe the inability of the system to adapt to the rapidly varying parameters when $T$ is small ($\sim 0.04\,\text{ms}$), as in Fig. 4d (labelled 5 and 6). The chiral behaviour originates from the asymmetry between the imaginary dynamical phase factors acquired during the counterclockwise and clockwise encirclings, which results in preferential amplification, leading to imbalanced spin populations. A more detailed classification of the behaviour of $C$ in different phases is provided in the Supplementary Information (Supplementary Fig. 4).

Our system provides an intriguing platform to study the interplay between many-body statistics and non-Hermiticity. However, the current observations can be effectively described within the non-Hermitian picture by ignoring a quantum jump operator in the Lindblad equation. This is an appropriate approximation for a large number of nearly non-interacting atoms, smearing out the effect of the quantum jumps24,42. Furthermore, hole excitations of the Fermi sea (which requires a full master equation approach) are rapidly relaxed in the timescale of the Fermi energy due to the decoherence heating process induced by the loss beam. This rapid relaxation allows us to treat the system as a $2 \times 2$ non-Hermitian effective Hamiltonian.

In this Letter we have experimentally demonstrated how non-Hermiticity affects the dispersion relation of spin–orbit-coupled fermions inducing the $PT$ symmetry-breaking transition. The topological nature near the EP indicates that non-Hermiticity can fundamentally modify the physical properties of the spin–orbit-coupled quantum system. In contrast to classical systems where only single (bosonic) particle dynamics are considered, our system sets the stage for investigating many-body interacting fermions with dissipation39. Furthermore, the possibility of exploring non-equilibrium dynamics, quantum thermodynamics40 and information criticality41 across the $PT$ symmetry-breaking transition by engineering the non-Hermitian Hamiltonian in a dynamic manner is conceivable.

Online content
Any methods, additional references, Nature Research reporting summaries, source data, extended data, supplementary information, acknowledgements, peer review information; details of author contributions and competing interests; and statements of data and code availability are available at https://doi.org/10.1038/s41567-021-01491-x.

Received: 29 September 2021; Accepted: 10 December 2021; Published online: 24 January 2022

References
1. Hasan, M. & Kane, C. Colloquium: topological insulators. Rev. Mod. Phys. 82, 3045 (2010).
2. Ashida, Y., Gong, Z. & Ueda, M. Non-Hermitian physics. Adv. Phys. 69, 249–435 (2020).
3. Miri, M.-A. & Alu, A. Exceptional points in optics and photonics. Science 363, eaar7709 (2019).
4. Xu, Y., Wang, S.-F. & Duan, L.-M. Weyl exceptional rings in a three-dimensional dissipative cold atomic gas. Phys. Rev. Lett. 118, 045701 (2017).
5. Gong, Z. et al. Topological phases of non-Hermitian systems. Phys. Rev. X 8, 031079 (2018).

Fig. 2 | Momentum-resolved Rabi spectroscopy for observing band merging. a, Single-particle energy dispersion of dressed states with increasing dissipation. The energy gap at the resonant quasi-momentum $q_0 = 1.0k_b$ is closed above the EP at $\gamma_r = 9.75E_r$. b, Rabi oscillation for $\delta = +4E_r$ and $\Omega_r = 4.5E_r$. Atoms, prepared in $|\uparrow\rangle$, are suddenly projected into a superposition of eigenstates revealing a non-uniform energy gap. The spin polarization is averaged taking into account the finite optical resolution over $\pm 0.15k_b$. Lines in b show a damped sinusoidal curve.
**Fig. 3 |** \( \mathcal{PT} \) symmetry-breaking transition and band closing at the EP. **a,** Time-dependent spin texture obtained for different dissipation strengths. The quasimomentum-resolved spin polarization is monitored after switching on SOC fields with dissipation for a variable time. Each spin texture is averaged over ten measurements. **b,** Energy bandgap (circles) measured via Rabi spectroscopy, which is in good agreement with theory (dashed lines). We estimate the uncertainty of the theoretical bandgap (shaded region) based on calibrated atom loss and its measurement uncertainty. **c,** The unbroken phase with prominent Rabi oscillations (left) and the broken phase showing a monotonic spin polarization (right). Lines in **c** show a fitted damped sinusoidal curve. **d-e,** Phase diagram of the \( \mathcal{PT} \) symmetry-breaking transition. The energy gap and damping rate measured from the Rabi oscillation are shown with the real (d) and imaginary (e) eigenvalues of \( \mathcal{H} \) (solid lines), respectively. The shaded region indicates the uncertainty of the bandgap and damping rate associated with the uncertainty of atom loss. The inset in **d** shows the bandgap measured at \( q_x = 0.5\kappa \). The error bars in all panels represent standard fitting errors (vertical) and the experimental uncertainty related to the calibration (horizontal).

6. Yao, S., Song, F. & Wang, Z. Non-Hermitian Chern bands. *Phys. Rev. Lett.* 121, 136802 (2018).
7. Helbig, T. et al. Generalized bulk-boundary correspondence in non-Hermitian topoelectrical circuits. *Nat. Phys.* 16, 747–750 (2020).
8. Xiao, L. et al. Non-Hermitian bulk-boundary correspondence in quantum dynamics. *Nat. Phys.* 16, 761–766 (2020).
9. Weidemann, S. et al. Topological funneling of light. *Science* 368, 311–314 (2020).
10. Lin, Y.-J., Jimenez-Garcia, K. & Spielman, I. B. Spin–orbit coupled Bose–Einstein condensates. *Nature* 471, 83–86 (2011).
11. Wang, P. et al. Spin–orbit coupled degenerate Fermi gases. *Phys. Rev. Lett.* 109, 095301 (2012).
12. Cheuk, L. W. et al. Spin-injection spectroscopy of a spin–orbit coupled Fermi gas. *Phys. Rev. Lett.* 109, 095302 (2012).
13. Huang, L. et al. Experimental realization of two-dimensional synthetic spin–orbit coupling in ultracold Fermi gases. *Nat. Phys.* 12, 540–544 (2016).
14. Wu, Z. et al. Realization of two-dimensional spin–orbit coupling for Bose–Einstein condensates. *Science* 354, 83–88 (2016).
15. Song, B. et al. Observation of symmetry-protected topological band with ultracold fermions. *Sci. Adv.* 4, eaao4748 (2018).
16. Song, B. et al. Observation of nodal-line semimetal with ultracold fermions in an optical lattice. *Nat. Phys.* 15, 911–916 (2019).
17. Wang, Z.-Y. et al. Realization of an ideal Weyl semimetal band in a quantum gas with 3D spin–orbit coupling. *Science* 372, 271–276 (2021).
18. Yamamoto, K. et al. Theory of non-Hermitian fermionic superfluidity with a complex-valued interaction. *Phys. Rev. Lett.* 123, 123601 (2019).
19. Lazarides, A., Roy, S., Piazza, F. & Moessner, R. Time crystallinity in dissipative Floquet systems. *Phys. Rev. Res.* 2, 022002 (2020).
20. Luo, X.-W. & Zhang, C. Higher-order topological corner states induced by gain and loss. *Phys. Rev. Lett.* 123, 073601 (2019).
21. Xu, H., Mason, D., Jiang, L. & Harris, J. G. E. Topological energy transfer in an optomechanical system with exceptional points. *Nature* 537, 80–83 (2016).
22. Doppler, J. et al. Dynamically encircling an exceptional point for asymmetric mode switching. *Nature* 537, 76–79 (2016).
23. Hassan, A. U. et al. Chiral state conversion without encircling an exceptional point. *Phys. Rev. A* 96, 052129 (2017).
24. Wiersig, J. Enhancing the sensitivity of frequency and energy splitting detection by using exceptional points: application to microcavity sensors for single-particle detection. *Phys. Rev. Lett.* 112, 203901 (2014).
25. Liebester, M. et al. Pump-induced exceptional points in lasers. *Phys. Rev. Lett.* 108, 173901 (2012).
Fig. 4 | Topological chiral spin transfer by dynamically encircling an EP. a. A Fermi gas is first adiabatically prepared in the initial $|\downarrow\rangle$-dominated state (spin polarization indicated by blue) in the parameter space with $\gamma_z=0$, $E_r$, and $\delta=3E_r$ (left). We effectively encircle the EP with fermions by tuning the loss and two-photon detuning (green arrow). The spin polarization before dynamic encircling is $-0.7$ in the initial state. The right panel of a shows the expected state conversion efficiency $C$ from $|\downarrow\rangle$ to $|\uparrow\rangle$ against loop duration time $T$ and the maximum atom loss $\gamma^\text{max}$ for counterclockwise (CCW, top) and clockwise (CW, bottom) encircling in the parameter space. The dominant spin polarization of the quantum state measured after the encircling is indicated by the colour of the circles and stars. The final spin polarization after encircling is experimentally measured for different $T$ and $\gamma^\text{max}$ as indicated by the colour of the circle or star, after following the trajectory along clockwise and counterclockwise closed loops. The horizontal and vertical dashed lines indicate the EP and $T=10.1\,\text{ms}$, respectively. b, c, Chiral behaviour of the final spin polarization for the experimental configurations indicated as circles with the indices as in a (1 and 2 in b, 3 and 4 in c). For other values of $T$, we keep the same fraction of time for each edge of the trajectory as in the experiment. d, No chirality is observed when dynamic encircling is instantaneous, within 0.04 ms, corresponding to the filled star in the right panel of a. For all cases in b–d, corresponding histograms for the 100 consecutive measurement series with a binning width of 0.02 for spin polarization are shown together with the encircling paths of the Fermi sea (green region) in the parameter space where the exceptional point is indicated by the yellow circle and $k=q_x-q_y/4$ (see Supplementary Information for more details). Error bars represent the standard deviation of the measurements.

26. Zhen, B. et al. Spawning rings of exceptional points out of Dirac cones. Nature 525, 354–358 (2015).
27. Zhao, H. et al. Non-Hermitian topological light steering. Science 365, 1163–1166 (2019).
28. Minganti, F., Miranowicz, A., Chhajlany, R. W. & Nori, F. Quantum exceptional points of non-Hermitian Hamiltonians and Liouvillians: the effects of quantum jumps. Phys. Rev. A 100, 062131 (2019).
29. Xiao, L. et al. Observation of topological edge states in parity-time-symmetric quantum walks. Nat. Phys. 13, 1117–1121 (2017).
30. Ozturk, F. E. et al. Observation of a non-Hermitian phase transition in an optical quantum gas. Science 372, 88–91 (2021).
31. Naghiloo, M., Abbasi, M., Joglekar, Y. N. & Murch, K. W. Quantum state tomography across the exceptional point in a single dissipative qubit. Nat. Phys. 15, 1232–1236 (2019).
32. Wu, Y. et al. Observation of parity-time symmetry breaking in a single-spin system. Science 364, 878–880 (2019).
33. Gao, T. et al. Observation of non-Hermitian degeneracies in a chaotic exciton-polariton billiard. Nature 526, 554–558 (2015).
34. Li, J. et al. Observation of parity-time symmetry breaking transitions in a dissipative Floquet system of ultracold atoms. Nat. Commun. 10, 855 (2019).
35. Takasu, Y. et al. PT-symmetric non-Hermitian quantum many-body system using ultracold atoms in an optical lattice with controlled dissipation. Prog. Theor. Exp. Phys. 2020, ptaa094 (2020).
36. Liu, W., Wu, Y., Duan, C.-K., Rong, X. & Du, J. Dynamically encircling an exceptional point in a real quantum system. Phys. Rev. Lett. 126, 170506 (2021).
37. Kofman, A. G. & Kurizki, G. Acceleration of quantum decay processes by frequent observations. Nature 405, 546–550 (2000).
38. Mailybaev, A. A., Kirillov, O. N. & Seyranian, A. P. Geometric phase around exceptional points. Phys. Rev. A 72, 014104 (2005).
39. Uzdin, R., Mailybaev, A. & Moiseyev, N. On the observability and asymmetry of adiabatic state flips generated by exceptional points. J. Phys. A Math. Theor. 44, 435302 (2011).
40. Wu, J. H., Artoni, M. & Rocca, G. C. L. Non-Hermitian degeneracies and unidirectional reflectionless atomic lattices. Phys. Rev. Lett. 113, 123004 (2014).
41. Yoon, J. W. et al. Time-asymmetric loop around an exceptional point over the full optical communications band. Nature 562, 86–90 (2018).
42. Daley, A. J. Adv. Phys. 63, 77 (2014).
43. Zhou, L., Yi, W. & Cui, X. Dissipation-facilitated molecules in a Fermi gas with non-Hermitian spin-orbit coupling. Phys. Rev. A 102, 043310 (2020).
44. Deffner, S. & Saxena, A. Jarynski equality in PT-symmetric quantum mechanics. Phys. Rev. Lett. 114, 150601 (2015).
45. Kawabata, K., Ashida, Y. & Ueda, M. Information retrieval and criticality in parity-time-symmetric systems. Phys. Rev. Lett. 119, 190401 (2017).
Methods
Preparation of the sample. We slowed down $^{173}$Yb atoms in a two-stage process and pre-cooled them in an intercombination magneto-optical trap. We then began with an experiment using a two-component degenerate $^{173}$Yb Fermi gas ($2 \times 10^5$ atoms) prepared at $T/T_F \leq 0.3(1)$, where $T_F = k_B \times 400$ nK, following forced evaporative cooling in an optical dipole trap formed by 1.064-nm and 532-nm laser beams with a trap frequency of $\omega = (\omega_x, \omega_y, \omega_z)^{1/2} = 112 \times 2\pi$ Hz. Here $|\uparrow, \downarrow\rangle = |m_x = 5/2, m_y = 3/2\rangle$ represent hyperfine states of the ground manifold with a negligible interaction strength of $k_d \approx 0.1$ where $k_d$ is the Fermi wavevector and $d$ denotes the $s$-wave scattering length. A quantized axis was fixed by a bias magnetic field of 13.6 G along the $z$ direction.

To create a non-Hermitian SOC Hamiltonian with equal strengths of the Rashba and Dresselhaus SOC fields, two hyperfine states of the ground-state manifold of $^{173}$Yb (labelled as the $|\uparrow\rangle$ and $|\downarrow\rangle$ states) were coupled by a pair of Raman transition beams intersecting at $\theta = 76^\circ$ and blue-detuned by $-1$ GHz. Typically, before the SOC was switched on by a pair of two-photon Raman beams, two hyperfine levels ($|\uparrow\rangle$, $|\downarrow\rangle$) were isolated from other states (that is, $m_x = 1/2$, $-1/2$, $-3/2$ and $-5/2$) using the spin-dependent light shift induced by a $\sigma^-$ polarized beam (referred to as a lift beam), blue-detuned by $-1$ GHz, which lifted the degeneracy of the ground manifold. Within experimental resolution, no atoms were observed in the hyperfine states other than $|\uparrow\rangle$ and $|\downarrow\rangle$. The background heating arising from the Raman transition was negligible within the timescale of the experiment.

Control of atom loss. To control the dissipation of the system, we used a dedicated plane-wave laser beam (referred to as the loss beam) at a small detuning around the 556-nm narrow intercombination transition $|3S_1/2, F = 2, m_F = 2\rangle \leftrightarrow |3P_1/2, F = 2, m_F = 2\rangle$ with the natural linewidth of 182 kHz. The loss beam is $\sigma^+$ polarized and is detuned by 1.2 MHz and 6.9 MHz with respect to the $|m_x = 3/2\rangle \rightarrow |m_y = 1/2\rangle$ and $|m_x = 5/2\rangle \rightarrow |m_y = 3/2\rangle$ transitions, respectively, which results in spin-sensitive photon scattering. The detuning of the loss beam is chosen such that the light is essentially resonant for atoms in the $|m_x = 3/2\rangle \rightarrow |\downarrow\rangle$ state with a fixed ratio $r_L/r_P$ of 13. Although the absorption and re-emission of a photon can leave an atom back in its original state, for example, with a probability of $\sim 14\%$ for atoms in $|\downarrow\rangle$ (see the description of the relative optical transition strength between hyperfine states in the Supplementary Information), we define the effective photon scattering rate for the genuine one-body dissipation, ignoring the heating effect, which is negligible in our current experimental regime. Nevertheless, the sample heats up within the timescale of $E_r^{-1}$ in the strong dissipation regime, for example, at a loss rate of $r_L = 10 E_r$, resulting in $T \approx 140$ nK. We expect that the heating associated with the photon scattering can be further suppressed by pumping excited atoms into the $|3P_1/2, m_F = 0\rangle$ manifold to $|3S_1/2, m_F = 0\rangle$ with 680-nm light. We calibrated the photon scattering rate $\gamma_p$ by fitting a function $e^{-\gamma_p t}$ to the atom number of the state $|\sigma = \uparrow, \downarrow\rangle$ after the loss beam was suddenly switched on (see Supplementary Information for more details).

We achieved a tunable loss rate by controlling the power of the loss beam. The change in the two-photon detuning ($\delta$) due to the energy shift induced by the loss beam was less than $\hbar \times 0.1$ kHz at $r_L = 1 E_r$.

$PT$-symmetric Hamiltonian. The Hamiltonian $\mathcal{H}$ can be decomposed as $\mathcal{H} = \mathcal{H}_{PT} - i\frac{\mathcal{D}_{PT}}{\hbar} I$, where we define the $PT$-symmetric Hamiltonian as

$$\mathcal{H}_{PT} = \sum_{\sigma = \uparrow, \downarrow} \left( \frac{\hbar^2}{2m} (q_x - k) \right)^2 |\uparrow\rangle \langle \uparrow| + \frac{\hbar^2}{2m} (q_y + k) \right)^2 |\downarrow\rangle \langle \downarrow| + \frac{\hbar^2}{2m} (q_z + k) \right)^2 |\uparrow\rangle \langle \downarrow| + \frac{\hbar^2}{2m} (q_z - k) \right)^2 |\downarrow\rangle \langle \uparrow|$$

$\mathcal{P} = \sigma$, represents the spin exchange operation and $T = \hbar$ denotes the pseudo time-reversal operation with $K$ the complex conjugate operation, which results in $[\mathcal{H}_{PT}, \mathcal{P}] = 0$ when $q_z = (\frac{\hbar}{\Omega}) k$. The constant decay term $\frac{\hbar^2}{\Omega^2} \frac{\mathcal{D}_{PT}}{\hbar} I$ does not affect the $PT$ symmetry-breaking transition. Alternatively, the gain and loss can be understood as being balanced at the EP. Near the gap where SOC is resonant (that is, $q_z = (\frac{\hbar}{\Omega}) k$), the dissipative two-level system is effectively described by $\mathcal{H}_{PT}$.

In the strong dissipation limit (that is, parity-time symmetry-broken phase), the slowly decaying eigenstate (that is, complex energy) has a near-unity overlap with the spin-up state without SOC, while the rapidly decaying eigenstate becomes similar to the bare spin-down state. When the dissipation is extremely strong, the quantum dynamics of the slowly decaying state is effectively projected onto the low-loss manifold, which is a bare energy band of spin-up atoms. Therefore, the SOC effectively disappears in this limit, closing the bandgap.

Data availability
The data that support the findings of this work are available from the corresponding authors upon reasonable request.

Acknowledgements
G.-B.J. acknowledges support from the RGC and the Croucher Foundation through grants nos. 16305317, 16304918, 16306119, 16302420, C 6005-17G and N-HKUST601/117. G.-B.J. is further supported by the Harilela foundation. I.L. acknowledges support from the RGC through grants nos. 16304520 and C6013-18G.

Author contributions
Z.R., E.Z., C.H. and K.K.P. carried out the experiment and data analysis and helped with numerical calculations. D.L. performed theoretical calculations. G.-B.J. and I.L. supervised the research.

Competing interests
The authors declare no competing interests.

Additional information
Supplementary information The online version contains supplementary material available at https://doi.org/10.1038/s41567-021-04919-x.

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Peer review information Nature Physics thanks Wei Yi and the other, anonymous, reviewer(s) for their contribution to the peer review of this work.

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