astrophysical simulations (Fig. 3A; see methods), performed with mass-ejection rates (22) and magnetic field strengths (10) typical of YSOs. We further verified that the range of density and temperature conditions found along the jet implies that the material within covers the full range of ionization fraction, with most of the jet part being only weakly ionized, consistent with observations (23). We see the principal dynamics and features of the laboratory flows also in the simulation results, which support the scalability of the laboratory experiment. We varied the wind and field parameters extensively with detailed numerical simulations to verify that the poloidal-field—induced jet collimation mechanism was a very robust process.

An interesting outcome of our study is that typical YSOs should then have a stationary region of shock-heated plasma that forms within a few tens to 100 AU from the wind source (see Fig. 3A). This numerical prediction can be directly compared to the analysis of observations made over more than one decade using x-ray satellites that have revealed (see Fig. 3C) bright sources of stationary x-ray emission zones located at the base of jets (~100 AU from the source) emerging from a YSO (24–26) and distinct from it. They have not been explained so far by self-collimation models of jet formation but are still consistent with the process revealed here (Fig. 3B).

We have therefore proposed a simple and plausible scenario for the collimation of a narrow stable jet past its launching phase (1, 2) that is consistent with recent astrophysical observations (20, 24, 25). In addition to helping to advance the understanding of jet—core interaction in YSOs, our work enables studying and/or modeling important aspects of jet physics in the laboratory. These include, e.g., transverse instabilities that can affect the jet structure, or episodic ejections, i.e., multicomponent and time-dependent interacting winds, which can be easily simulated in the laboratory by using multiple laser pulses separated by a few nanoseconds. Producing such magnetized narrow plasma columns and letting them strike a solid will also uniquely allow the study of plasma dynamics in accretion columns in young stars; that is, one can model magnetic arches that are loaded with disk material that free-falls toward the star. Beyond these aspects, adapting the present experimental work to other configurations will permit advances in resolving pending questions but in a wide range of astrophysical and plasma physics systems where magnetic fields are thought to play an important role.

REFERENCES AND NOTES

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SUPPLEMENTARY MATERIALS
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OPTICS
Loss-induced suppression and revival of lasing
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Controlling and reversing the effects of loss are major challenges in optical systems. For lasers, losses need to be overcome by a sufficient amount of gain to reach the lasing threshold. In this work, we show how to turn losses into gain by steering the parameters of a system to the vicinity of an exceptional point (EP), which occurs when the eigenvalues and the corresponding eigenstates of a system coalesce. In our system of coupled microresonators, EPs are manifested as the loss-induced suppression and revival of lasing. Below a critical value, adding loss annihilates an existing Raman laser. Beyond this critical threshold, lasing recovers despite the increasing loss, in stark contrast to what would be expected from conventional laser theory. Our results exemplify the counterintuitive features of EPs and present an innovative method for reversing the effect of loss.

Dissipation is ubiquitous in nature; the states of essentially all physical systems thus have a finite decay time. A proper description of this situation requires a departure from conventional Hermitian models with real eigenvalues and orthogonal eigenstates to non-Hermitian models featuring complex eigenvalues and nonorthogonal eigenstates (1–7). When tuning the parameters of such a dissipative system, its complex eigenvalues and the corresponding eigenstates may coalesce, giving rise to a non-Hermitian degeneracy, also called an exceptional point (EP) (4). The presence of such an EP has a dramatic effect on the system, leading to nontrivial physics with interesting counterintuitive features such as resonance trapping (5), a mode exchange when encircling an EP (6), and a singular topology in the parameter landscape (7). These characteristics can control the flow of light in optical devices with both loss and gain. In particular, waveguides with parity-time symmetry (8), where loss and gain are balanced, have attracted enormous attention (9, 10), with effects
such as loss-induced transparency (11), unidirectional invisibility (12), and reflectionless scattering (13, 14) having already been observed.

Theoretical work indicates that EPs give rise to many more intriguing effects when they occur near the lasing regime; for example, enhancement of the laser linewidth (15, 16), fast self-pulsations (15), and a pump-induced lasing death (17). Realizing such anomalous phenomena, however, requires moving from waveguides to resonators, which can trap and amplify light resonantly beyond the lasing threshold. With the availability of such devices (18, 19), we discuss here the most counterintuitive aspect, namely, that close to an exceptional point, lasing should be inducible solely by adding loss to a resonator.

Our experimental system (20) consists of two directly coupled silica whispering-gallery-mode resonators (WGMRs), $\mu R_1$ and $\mu R_2$, each coupled to a different fiber-taper, WG1 and WG2 (Fig. 1A and supplementary text S1). The resonance frequencies of the WGMRs were tuned to be the same via the thermo-optic effect, and a controllable coupling strength $k$ was achieved between the WGMRs by adjusting the interresonator distance. To observe its behavior in the vicinity of an EP, the system was steered parametrically via $k$ and an additional loss $\gamma_{\text{tip}}$ induced on $\mu R_2$ by a chromium (Cr)–coated silica-nanofiber tip (Fig. 1, B and C), which features strong absorption in the 1550-nm band. The strength of $\gamma_{\text{tip}}$ was increased by enlarging the volume of the nanotip within the $\mu R_2$ mode field, resulting in a broadened resonance linewidth with no observable change in resonance frequency (Fig. 1D). A small fraction of the scattered light from the nanotip coupled back into $\mu R_2$ in the counter-propagating (backward) direction and led to a resonance peak whose linewidth broadened as the loss increased (Fig. 1E). The resonance peak in the backward direction was $\sim 10^4$ of the input field, confirming that the linewidth broadening and the decrease of the resonance depth in the forward direction were due to $\gamma_{\text{tip}}$ via absorption and scattering losses, but not due to backscattering into the resonator.

In the first set of experiments, WG2 was moved away from $\mu R_2$ to eliminate the coupling between them. We investigated the evolution of the eigenfrequencies and the transmission spectra $T_{1 \rightarrow 2}$ from input port 1 to output port 2 by continuously increasing $\gamma_{\text{tip}}$ while keeping $k$ fixed. In this configuration, losses experienced by $\mu R_1$ and $\mu R_2$ were $\gamma_1' = \gamma_1 + \gamma_{\text{tip}}$ and $\gamma_2' = \gamma_2 + \gamma_{\text{tip}}$, where $\gamma_1$ is the WG1-$\mu R_1$ coupling loss, and $\gamma_1$ and $\gamma_2$ include material absorption, scattering, and radiation losses of $\mu R_1$ and $\mu R_2$. The coupling between the WGMRs led to the formation of two supermodes with complex eigenfrequencies $\omega_1 = \omega_0 - ik_2 + \beta$ whose real and imaginary parts are respectively denoted by $\omega_1'$ and $\omega_1''$ (20). Here, $\omega_0$ is the resonance frequency of the solitary WGMRs; $\chi = (\gamma_1' + \gamma_2')/4$ and $\Gamma = (\gamma_1' - \gamma_2')/4$, respectively, quantify the total loss and the loss contrast of the WGMRs; and $\beta = \sqrt{\chi^2 - \Gamma^2}$ reflects the transition between the strong and the weak intermode coupling regimes due to an interplay of the interresonator coupling strength $k$ and the loss contrast $\Gamma$ (supplementary text S1). In the strong-coupling regime, quantified by $k > |\Gamma|$, and real $\beta$, the supermodes had different resonance frequencies (mode splitting of 2β) but the same linewidths quantified by $\chi$. This was reflected as two spectrally separated resonance modes in $T_{1 \rightarrow 2}$ (Fig. 2A (i)) and in the corresponding eigenfrequencies (Fig. 2B (i)). Because our system satisfied $\gamma_1' + \gamma_2' > \gamma_2$, introducing $\gamma_{\text{tip}}$ to $\mu R_2$ increased the amount of splitting until $\gamma_1' + \gamma_{\text{tip}} = (\gamma_2 + \gamma_{\text{tip}})$ (that is, $\gamma_1' = \gamma_2'$) was satisfied (Fig. 2, A (ii) and B (ii)). Increasing $\gamma_{\text{tip}}$ beyond this point gradually led to an overlap of the supermode resonances (Fig. 2A (iii)), such as to necessitate a fit to a theoretical model to extract the complex resonance parameters (supplementary text S1 and S2) (20). At $\gamma_{\text{tip}} = \gamma_{\text{tip}}^* \gamma_2$, the supermodes coalesced at the EP. With a further increase of $\gamma_{\text{tip}}$, the system entered the weak-coupling regime, quantified by $k < |\Gamma|$ and imaginary $\beta$, which led to two supermodes with the same resonance frequency but different linewidths (Fig. 2, A (iv) and B (iv)). The resulting resonance trajectories in the complex plane clearly displayed a reversal of eigenvalue evolution (Fig. 2B): The real parts of the eigenfrequencies of the system approached each other while their imaginary parts remained equal until the EP. After passing the EP, their imaginary parts were repelled, which resulted in an increasing imaginary part for one of the eigenfrequencies and a decreasing imaginary part for the other. As a result, one of the modes became less lossy, whereas the other became more lossy (supplementary text S4).

By repeating the experiments for different $k$ and $\gamma_{\text{tip}}$, we obtained the eigenfrequency surfaces $\omega_0(k, \gamma_1')$, whose real and imaginary parts are shown in Fig. 2, C and D. The resulting surfaces exhibit a complex square root–function topology with the special feature that a coalescence of the eigenfrequencies can be realized by varying either $k$ or $\gamma_{\text{tip}}$, alone, which leads to a continuous thread of EPs along what may be
called an exceptional line. As expected, the slope of this line is such that stronger $k$ requires higher $\gamma_{\text{tip}}$ to reach the EP (supplementary text S4).

Our second set of experiments was designed to elucidate the effect of the EP on the intracavity field intensities. For this we used both WGI and WG2, introducing an additional coupling loss $g_{c2}$ to $m_{R2}$ (that is, $\gamma_{c2} = \gamma_{2} + \gamma_{\text{tip}} + \gamma_{c2}$). We tested two different cases by choosing different mode

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Fig. 2. Evolution of the transmission spectra and the eigenfrequencies as a function of loss $\gamma_{\text{tip}}$ and interresonator coupling strength $k$. (A) Transmission spectra $T_{1 \rightarrow 2}$ showing the effect of loss on the supermodes. Blue and red curves denote the experimental data and the best fit using a theoretical model (supplementary text S1 and S2), respectively. (B) Evolution of the eigenfrequencies of the supermodes in the complex plane as $\gamma_{\text{tip}}$ was increased. Open circles and squares are the eigenfrequencies estimated from the measured $T_{1 \rightarrow 2}$. Dashed red and blue lines denote the best theoretical fit to the experimental data. (C and D) Eigenfrequency surfaces in the ($k, \gamma_{c2}$) parameter space (supplementary text S4).

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Fig. 3. Loss-induced enhancement of intracavity field intensities and thermal nonlinearity in the vicinity of an exceptional point. (A and B) Intracavity field intensities of the resonators at $\omega_1$ (blue, $I_1$ of $\mu_{R1}$; green, $I_2$ of $\mu_{R2}$; red, total $I$) for (A) case 1 and (B) case 2. Normalization was done with respect to the total intensity at $\gamma_{\text{tip}} = 0$. a.u., arbitrary units. (C) Total intracavity field intensities $I_T$ at eigenfrequencies $\omega_{\pm}$ (black) and $\omega_0$ (red) for case 1 (see supplementary text S5 for case 2). Normalization was done with respect to the intensity at the EP. (D) Effect of loss on nonlinear thermal response of coupled resonators (supplementary text S8): (i) solitary resonator, (ii) coupled resonators with $\gamma_{\text{tip}} = 0$, and (iii and iv) coupled resonators with increasing $\gamma_{\text{tip}}$ (20). Circles in (A) to (C) and squares in (C) were calculated from experimentally obtained transmissions $T_{1 \rightarrow 2}$ and $T_{1 \rightarrow 4}$, whereas solid and dashed curves are from the theoretical model (supplementary text S1 to S3).
pairs in the resonators (20). In case 1, the mode in \( \mu_{R1} \) had a higher loss than the mode in \( \mu_{R2} \) \((\gamma_1 + \gamma_0 > \gamma_2 + \gamma_0)\); in case 2, the mode in \( \mu_{R2} \) had a higher loss \((\gamma_1 + \gamma_0 < \gamma_2 + \gamma_0)\). The system was adjusted so that two spectrally separated supermodes were observed in the transmission spectra \( T_{1 \rightarrow 2} \) and \( T_{1 \rightarrow 4} \) as resonance dipped and peaked (supplementary text S3). No resonance dip or peak was observed at port 3. Using experimentally obtained \( T_{1 \rightarrow 2} \) and \( T_{1 \rightarrow 4} \), we estimated the intracavity fields \( I_1 \) and \( I_2 \) and the total intensity \( I_T = I_1 + I_2 \) as a function of \( \gamma_{\text{tip}} \) (Fig. 3, A to C, and supplementary text S1, S2, S5, and S6). As \( \gamma_{\text{tip}} \) was increased, \( I_T \) first decreased and then started to increase despite increasing loss. This loss-induced recovery of the intensity is in contrast to the expectation that the intensity would decrease with increasing loss and is a direct manifestation of the EP.

The effect of increasing \( \gamma_{\text{tip}} \) on \( I_1 \) and \( I_2 \) at \( \omega_0 \) is depicted in Fig. 3, A and B. When \( \gamma_{\text{tip}} = 0 \) and the system was set in the strong-coupling regime, the light input at \( \mu_{R1} \) was freely exchanged between the resonators, establishing evenly distributed supermodes. As a result, the intracavity field intensities were almost equal. As \( \gamma_{\text{tip}} \) was increased, \( I_1 \) and \( I_2 \) decreased continuously at different rates until \( I_1 \) reached a minimum at \( \gamma_{\text{tip}} = \gamma_{\text{tip}}^{\text{min}} \). The rate of decrease was higher for \( I_2 \) because of increasingly higher loss of \( \mu_{R2} \). Beyond \( \gamma_{\text{tip}}^{\text{min}} \), until the EP was reached at \( \gamma_{\text{tip}} = \gamma_{\text{EP}} \), the system remained in the strong-coupling regime, but the supermode distributions were strongly affected by \( \gamma_{\text{tip}} \), which led to an increase of \( I_1 \) and, hence, of \( I_T \), whereas no appreciable change was observed for \( I_2 \) (20). Increasing \( \gamma_{\text{tip}} \) further pushed the system beyond the EP, thereby completing the transition from the strong-coupling to the weak-coupling regime during which \( I_T \) increased and kept increasing, whereas \( I_2 \) continued decreasing. This behavior is a manifestation of the progressive localization of one of the supermodes in the less lossy \( \mu_{R2} \) and of the other supermode in the more lossy \( \mu_{R1} \). We conclude that the nonmonotonic evolution of \( I_T \) for increasing values of \( \gamma_{\text{tip}} \) is the result of a transition from a symmetric to an asymmetric distribution of the supermodes in the two resonators (supplementary text S5 and S7).

The data shown in Fig. 3, A and B, also demonstrate that the initial loss contrast of the resonators affects both the amount of \( \gamma_{\text{tip}} \) required to bring the system to the EP and the intensity values themselves (20): Increasing \( \gamma_{\text{tip}} \) in case 2 increased \( I_T \) to a higher value than that at \( \gamma_{\text{tip}} = 0 \); in case 1, on the other hand, \( I_T \) stayed below its initial value at \( \gamma_{\text{tip}} = 0 \). Finally, Fig. 3C shows that the intracavity field intensities at \( \omega_0 \) and \( \omega_0 \) coincided when \( \gamma_{\text{tip}} \geq \gamma_{\text{tip}}^{\text{EP}} \), that is, after the EP transition to the weak coupling regime (20). This is a direct consequence of the coalescence of eigenfrequencies \( \omega_1 \) at \( \omega_0 \) (supplementary text S5).

Whispering-gallery-mode microresonators (21) combine high-quality factor \( Q \) (long photon storage time, narrow linewidth) and high-finesse \( F \) (strong resonant power build-up) with microscale mode volume \( V \) (tight spatial confinement, enhanced resonant field intensity) and are thus ideal for studying quantum electrodynamics (22), optomechanics (23), lasing (24), and sensing (25–27). The ability of WGMRs to provide high intracavity field intensity and long interaction time reduces thresholds for nonlinear processes. Therefore, loss-induced reduction and the recovery of intracavity field intensities should have a direct effect on the thermal nonlinearity (28) and the Raman lasing (24, 29) in WGMRs.

Thermal nonlinearity in WGMRs is due to the temperature-dependent resonance-frequency shifts caused by material absorption of the intracavity field and the resultant heating (20, 28). In silica WGMRs, this is manifested as thermal broadening (or narrowing) of the resonance line when the wavelength of a probe laser is scanned from shorter to longer (or longer to shorter) wavelengths. In our system, thermal nonlinearity was observed in \( T_{1 \rightarrow 2} \) as a stark-fine feature (Fig. 3D). With an input power of 600 \( \mu \text{W} \), thermal broadening kicked in and made it impossible to resolve the individual supermodes (Fig. 3D (i and ii)). When \( \gamma_{\text{tip}} \) was introduced and gradually increased, thermal nonlinearity and the associated linewidth broadening decreased at first but then gradually recovered (Fig. 3D (iii and iv)). This aligns well with the evolution of the total intracavity field as a function of loss (supplementary text S8).

Finally, we tested the effect of the loss-induced recovery of the intracavity field intensity on Raman lasing in silica microtoroids (29, 30). The threshold for Raman lasing scales as \( P_{\text{laser}}^{\text{threshold}} = \frac{V}{g_T Q^2} \) (where \( g_T \) is the Raman gain coefficient), implying the importance of the pump intracavity field intensity in the process. With a pump in the 1550-nm band, Raman lasing in the silica WGMR takes place in the 1650-nm band. Figure 4 depicts the spectrum and the efficiency of Raman lasing in our system. The lasing threshold for \( \mu_{R2} \) was 150 \( \mu \text{W} \) (Fig. 4B, blue curve). Keeping the pump power fixed, we introduced \( \mu_{R1} \), which had a much larger loss than \( \mu_{R2} \). This effectively increased the total loss of the system and annihilated the laser (Fig. 4A, gray curve). Introducing \( \gamma_{\text{tip}} \) to \( \mu_{R2} \) helped to recover the Raman laser, whose intensity increased with increasing \( \gamma_{\text{tip}} \) (Fig. 4A). We also checked the lasing threshold of
each of the cases depicted in Fig. 4A and observed that as \( T_{\text{Ram}} \) was increased, the \( \Delta P_{\text{Ram}} \)-threshold increased at first but then decreased (Fig. 4B).

These observations are in stark contrast with what one would expect in conventional systems, where the higher the loss, the higher the lasing threshold. Surprisingly, in the vicinity of an EP, less loss is detrimental and annihilates the process of interest; more loss is good because it helps to recover the process. This counterintuitive effect happens because the supermodes of the coupled system readjust themselves as loss is gradually increased. When the loss exceeds a critical value, one supermode is mostly located in the subsystem with less loss and, thus, the total field can build up more strongly (20). As our results demonstrate, this behavior also affects nonlinear processes, such as thermal broadening and Raman lasing, that rely on intracavity field intensity.

Our system provides a comprehensive platform for further studies of EPs and opens up new avenues of research on non-Hermitian systems and their behavior. Our findings may also lead to new schemes and techniques for controlling and reversing the effects of loss in other physical systems, such as in photonic crystal cavities, plasmonic structures, and metamaterials.

REFERENCES AND NOTES
20. Supplementary materials are available on Science Online.

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QUANTUM ELECTRONICS

Cavity quantum electrodynamics with many-body states of a two-dimensional electron gas

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Light-matter interaction has played a central role in understanding as well as engineering new states of matter. Reversible coupling of excitons and photons enabled groundbreaking results in condensation and superfluidity of nonequilibrium quasiparticles with a photonic component. We investigated such cavity-polaritons in the presence of a high-mobility two-dimensional electron gas, exhibiting strongly correlated phases. When the cavity was resonant with the Fermi level, we observed previously unknown many-body physics associated with a dynamical hole-scattering potential. In finite magnetic fields, polaritons show distinct signatures of integer and fractional quantum Hall ground states. Our results lay the groundwork for probing nonequilibrium dynamics of quantum Hall states and exploiting the electron density dependence of polariton splitting so as to obtain ultrastrong optical nonlinearities.

Strong electric-dipole coupling of excitons in an intrinsic semiconductor quantum well (QW) to microcavity photons leads to the formation of mixed light-matter quasiparticles called cavity-polaritons (1, 2). In contrast to excitons, optical excitations from a two-dimensional electron gas (2DEG) show many-body effects associated with the abrupt turn-on of a valence-band hole-scattering potential. The resulting absorption spectrum exhibits a power-law divergence, stemming from final-state interaction effects that lead to orthogonal initial and final many-body wave functions, and is referred to as Fermi-edge singularity (3). Despite initial attempts (4–6), the nature of strong coupling of a Fermi-edge singularity, rather than an excitation, to a many-body mode remained largely unexplored.

We investigated a 2DEG in a modulation-doped GaAs QW embedded in a distributed Bragg reflector (DBR) microcavity (Fig. 1A). We used this system to analyze two distinct problems: In the absence of an external magnetic field (\( B_z = 0 \)), optical excitations from a 2DEG are simultaneously subject to Coulomb interactions and strong cavity coupling, which lead to the formation of Fermi-edge polaritons—a many-body excitation with a dynamical valence-hole scattering potential (7). In the \( B_z \neq 0 \) limit, we establish cavity quantum electrodynamics (QED) as a new paradigm for investigating integer and fractional quantum Hall (QH) states (8) in which information about a strongly correlated electronic ground state is transcribed onto the reflection and transmission amplitudes of a single-cavity photon.

The dispersion of the relevant electronic states when \( B_z = 0 \) is shown in Fig. 1B, left. The Fermi energy \( E_F \) is tuned by applying a gate voltage (\( V_G \)) between a p-doped GaAs top layer and the 2DEG. When \( B_z \neq 0 \), the electronic states form discrete Landau levels (LLs) (Fig. 1B, right). Each (spin-resolved) LL has a degeneracy of \( B_z/\hbar \); the filling factor \( \nu = n_e/\hbar (e/\nu) \) determines the number of filled LLs at a given electron density \( n_e \). We studied two different samples with gate-tunable \( \nu \) : sample \( A \) exhibits an electron mobility of \( \mu = 2 \times 10^5 \text{cm}^2/(\text{V s})^{-1} \) at \( n_e = 1.0 \times 10^{11} \text{cm}^{-2} \) (9). Sample \( B \) exhibits an order-of-magnitude-higher mobility of \( \mu = 2.5 \times 10^6 \text{cm}^2/(\text{V s})^{-1} \) at \( n_e = 2.3 \times 10^{11} \text{cm}^{-2} \).

The differential reflectivity (dR) spectra of sample \( A \) (blue) is shown in Fig. 1D as the electron density \( n_e \) in the 2DEG is varied from \( n_e < 2 \times 10^{10} \text{cm}^{-2} \) to \( n_e = 8 \times 10^{10} \text{cm}^{-2} \). The cavity mode is visible as a Dr peak at 1521 meV. Changing \( n_e \) modifies the nature of elementary optical excitations: At low electron densities

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