Wavelength-scale lasers provide promising applications through low power consumption requiring for optical cavities with increased quality factors. Cavity radiative losses can be suppressed strongly in the regime of optical bound states in the continuum; however, a finite size of the resonator limits the performance of bound states in the continuum as cavity modes for active nanophotonic devices. Here, we employ the concept of a supercavity mode created by merging symmetry-protected and accidental bound states in the continuum in the momentum space, and realize an efficient laser based on a finite-size cavity with a small footprint. We trace the evolution of lasing properties before and after the merging point by varying the lattice spacing, and we reveal this laser demonstrates the significantly reduced threshold, substantially increased quality factor, and shrunken far-field images. Our results provide a route for nanolasers with reduced out-of-plane losses in finite-size active nanodevices and improved lasing characteristics.
Ultra-compact lasers operating in the single-mode regime have been a long-standing goal for nanophotonics. Various mechanisms for strong light confinement were proposed and demonstrated in subwavelength and wavelength-scale optical cavities used to decrease their lasing threshold\textsuperscript{1-3}. However, nanolasers such as defect-type photonic crystal lasers and plasmonic lasers possess limited output power and exhibit instability to the structural disorder\textsuperscript{5,6}. These properties hinder the practical applications of nanolasers although they have small mode volumes and reduced laser thresholds. On contrary, a band-edge type laser with a periodic structure without defects has a relatively high threshold, but a high output power with the possibility of topological robustness\textsuperscript{7}.

Optical bound states in the continuum (BICs) were shown to be a versatile tool for substantial suppression of out-of-plane radiative losses and dramatic enhancement of the quality factor (Q factor) in infinite periodic structures to provide low-threshold lasing and high-output powers\textsuperscript{8-10}. The intrinsic topological nature of BICs splits them into a few groups, with two of the most conventional kinds represented by symmetry-protected BICs, existing at high-symmetry points of the momentum space, and accidental BICs, which can be realized for an arbitrary in-plane wavevector. Thus far, symmetry-protected BICs and accidental BICs were observed successfully in Si\textsubscript{3}N\textsubscript{4} and Si photonic crystal slabs\textsuperscript{11,12}. In addition, the lasing action was achieved for BIC cavities in a few recent experiments\textsuperscript{13-19}. The BIC-enhanced feature of a reasonably low threshold was successfully demonstrated in a cavity with lattices of finite size\textsuperscript{13}. However, most of the developed lasers were still based on cavities of substantially large scales with hundreds of periods and despite that, they did not demonstrate high Q factors. Very recently, a new kind of BIC mode was suggested, which we term here as a super-BIC (BIC in finite sizes). The design of the laser cavity is based on symmetry-protected and accidental BICs, which are dominated by the out-of-plane magnetic dipole, and agrees well with the fundamental nature of the TE-like mode (Supplementary Fig. 3).

Results

Infinite-size BIC cavity. The design of the laser cavity is based on an infinite-size InGaAsP photonic crystal slab structure with a thickness of 650 nm (inset in Fig. 1a) modulated with a square-lattice array of air holes. The calculated band structure shows that only the fundamental transverse-electric-like (TE-like) band is located within the emission wavelength range of the InGaAsP, ~1.5 μm, due to the thick slab of high-index material (Supplementary Fig. 1). The calculated map of Q factor in the k-space shows that the fundamental TE-like band demonstrates high Q factors at the origin and at specific points along highly symmetric Γ-X and Γ-M directions (Supplementary Fig. 2): in total, one symmetry-protected BIC (|k|= 0) and eight accidental BICs (|k| – k\textsubscript{c} = 0.067a/2n) are formed\textsuperscript{16}. In addition, the multipolar decomposition of electromagnetic fields inside the unit cell shows that both symmetry-protected and accidental BICs are dominated by the out-of-plane magnetic dipole, which agrees well with the fundamental nature of the TE-like mode (Supplementary Fig. 3). Each of accidental BICs is formed because of destructive interference between the dominating out-of-plane magnetic dipole and weak in-plane electric dipole, in-plane electric quadrupole, out-of-plane magnetic quadrupole. We also calculate the evolution of the polarization phases in the far-field for symmetry-protected and accidental BICs (Supplementary Fig. 4).\textsuperscript{26}

The calculation shows the topological charge q = 1 for the symmetry-protected BIC and four accidental BICs located on the Γ-M band, and q = -1 for four accidental BICs located on the Γ-X band. In addition, we calculate the wavelengths of the resonant modes shown in Fig. 1a. A symmetry-protected BIC and an accidental BIC mode appear in the range of the lattice constants a from 560 to 590 nm. The wavelengths of these modes become closer as the lattice constant increases at a < 576.3 nm, and only a single mode is observed after the two BIC modes merge at a = 576.3 nm.

Next, the normalized radiation loss of the fundamental TE-like band is calculated for different lattice constants along the Γ-X direction (black dots in Fig. 1b). At a = 568 nm, the zero-radiation loss is observed at k = 0 (symmetry-protected BIC) and at k = k\textsubscript{c} (accidental BIC), showing each sharp dependence on the wavevector. The position of the accidental BIC approaches k = 0 as the lattice constant increases; the merging of the BICs occurs at a = 576.3 nm and modifies the dependence on the wavevector in the vicinity of the band edge. The k-vector dependence returns to the narrow one at a = 590 nm when the single symmetry-protected BIC is restored. The radiation loss follows the laws |k(k – k\textsubscript{c})(k + k\textsubscript{c})|^2, k\textsuperscript{6}, or k\textsuperscript{2} depending on the lattice constant (red curves in Fig. 1b), because of the interaction of the topological charges q\textsubscript{a} at the symmetry-protected BIC and q\textsubscript{a} at the accidental BIC\textsuperscript{20}. This feature is also seen in the graph of the k vs. Q factor (Supplementary Fig. 1b).

Finite-size BIC cavity. The fast sixth-order dependence of the radiation losses on the wavevector is crucial to keep a high Q factor for the BIC mode in a finite-size cavity\textsuperscript{20}. Fig. 2a shows the mode magnetic field profile H\textsubscript{z} at the Γ point of the fundamental TE-like band in a photonic crystal slab with 15 × 15 periods. Unlike the infinite-size cavity, the mode profile shows an envelope distribution with a convex shape accompanied by mode leakage into free space\textsuperscript{25,27}. This finite-size effect results in the mode broadening Δk (white circle) in the k-space field distribution, FT(H\textsubscript{z}) (Fig. 2b), where FT means the spatial Fourier transformation. The mode broadening leads to increased radiation loss due to mixing with the off-Γ point modes with a finite Q factor. To increase the Q factor, it is essential to reduce the radiation loss due to the broadening (Fig. 2c). The undesired
radiation loss at off-Γ points can be effectively suppressed by moving the off-Γ BIC with the charge \( q_t \) inside the mode broadening range \( \Delta k \) (pre-merging regime) or to the Γ point \( k = 0 \) (merging regime). Notably, in the merging process with a limited number of air holes, the most effective radiation suppression can be achieved by placing the charge \( q_t \) at an optimum \( k_t \) in the pre-merging configuration rather than at \( k = 0 \) (Fig. 2c), which is explained below.

To elucidate the radiation loss mechanism in the finite-size system, we calculate the 2D maps of the radiation factor in momentum space with varying lattice constants (Fig. 2d). The radiation factor is defined by the \( k \)-space mode distribution in a finite domain and the Q factor in an infinite domain, i.e., \( |\text{FT}(H_\gamma)(\mathbf{k})/Q(\mathbf{k})| \), to account for both effects of the cavity size and the radiative loss. For example, the radiation factors are high before merging \((a = 568 \text{ nm})\) and after merging \((a = 578 \text{ nm})\) because of the substantial radiative loss at off-Γ points. At the exact merging condition of \( a = 576 \text{ nm} \), radiation loss still occurs at the boundary of the mode broadening. Interestingly, a much lower radiation factor is observed in the pre-merging regime at \( a = 573 \text{ nm} \) because the radiative loss is strongly suppressed in the large \( k \)-space area covered by \( q_t \) and \( q_\theta \). For more detailed quantitative analysis, we plot the inverse radiation factor, \( \Sigma |\text{FT}(H_\gamma)(\mathbf{k})/Q(\mathbf{k})|^{-1} \), as a function of the lattice constant when the hole diameter and slab thickness are fixed to 400 and 650 nm, respectively (Fig. 2e). The broad merging configuration occurs due to the finite-size effect. The graph becomes narrower and the maximum occurs at a larger lattice constant when the cavity size increases from \( N = 15 \) to 21. Further discussion is performed to investigate the influences caused by other structural parameters on the radiation factor (Supplementary Fig. 5).

In addition, we calculate the radiative Q factor using a full-wave numerical simulation (Fig. 2f). The agreement between the radiative Q factor and the inverse radiation factor demonstrates the effectiveness of our analysis based on the radiation factor. Consequently, in the finite-size cavity, the radiative loss near the merging-BIC regime (from pre-merging to merging) is still low, whereas the loss in the other BIC regimes is relatively high. Also, the optimal point with the lowest radiation in the finite and infinite-size cavities could be different from each other (Supplementary Note 1).

**Measurements of BIC lasing.** To experimentally verify the merging of the BICs, we fabricate square-lattice photonic crystal structures using a 650 nm-thick InGaAsP slab incorporating seven quantum wells (see “Methods” section). A set of samples with the lattice constant varying from 560 to 580 nm in 1 nm steps is fabricated with the hole radius fixed at \( \sim 200 \text{ nm} \) (Supplementary Fig. 6). The scanning electron microscope (SEM) images of the fabricated sample are shown in Fig. 3a. The photoluminescence (PL) measurements are performed using a 980-nm pulsed pump laser with a spot size of \( \sim 5.4 \mu \text{m} \) at room temperature (see “Methods” section).

One or two types of lasing modes are observed in the photonic crystal structures, depending on the lattice constant. The
measurement results of the lasing wavelength vs. the lattice constant in a sample are plotted in Fig. 3b. Single lasing peaks are observed at \(a > 568\) nm (bottom inset in Fig. 3b) and two peaks at \(a \leq 568\) nm. The wavelengths of the two peaks increase with different slopes as the lattice constant increases at \(a \leq 568\) nm. To identify these lasing modes, we measure their far-field images (see "Methods" section). Two distinct far-field images are observed, one with a highly confined donut shape (right inset; Fig. 3b) and the other with a widespread shape (left inset; Fig. 3b). Consequently, the lasing peaks can be classified into the two groups of symmetry-protected BIC (black dots) and accidental BIC modes (red dots), based on the measured far-field images and a comparison with the simulation results including those presented in Fig. 1a. We note that the merging occurs at \(a \approx 574\) nm by extrapolating the wavelength of the accidental BIC mode. Similar features are exhibited by the other samples with slightly different structural parameters (Supplementary Fig. 7).

To further investigate the optical properties of the symmetry-protected BIC laser, we compare the measured far-field images before merging (\(a = 568\) and 571 nm) with those after merging (\(a = 574\) and 578 nm) (Fig. 3c). The shapes of the lasing modes are identical, whereas the mode size decreases as the lattice constant increases and remains unchanged after merging. For more detailed quantitative analysis, we estimate the angle from the

![Graph showing measurement results of the lasing wavelength vs. lattice constant.](https://example.com/graph.png)
center to the first intensity maximum in eight measured far-field images (see “Methods” section). The change in the mode size is clearly shown in Fig. 3e (black dots). The angle decreases from ~4.9° to ~3.7° until \( a \geq 574 \) nm and remains almost constant at \( a \geq 574 \) nm. In fact, the size of the far-field image depends on the mode size in the near field\(^2\). Thus, by comparing the sizes of the measured and simulated far-field images, one can estimate the effective number of air holes (\( N \)) in the photonic crystal structure for the excitation of the corresponding near-field image. Our numerical simulations show that the structures with \( N = 19, 23, 27 \) and 27 yields the measured far-field images in Fig. 3c (left to right in Fig. 3d). A more systematic comparison between the measurement and simulation is shown in Fig. 3e, where the red dashed lines indicate the simulation results. Therefore, in the evolution from before-merging to after-merging, we observe that the after-merging BIC mode is lasing with a larger effective \( N \).

We term this BIC mode at \( a \geq 574 \) nm the super-BIC mode, as it possesses optical characteristics such as a single lasing peak and a shrunken far-field image that are distinct from those of the conventional BICs at \( a < 574 \) nm. In particular, the super-BIC mode is confined more strongly by the effectively increased number of air holes. This unique feature is useful for improving the laser performance despite the effect of the finite size on the Q factor.

**Lasing properties.** We examine the lasing properties of all demonstrated BIC lasers. First, the lasing spectra and light in–light out (L–L) curves are measured in the accidental BIC lasers (Supplementary Fig. 8); their threshold values are much larger than those of the symmetry-protected BIC lasers (Supplementary Fig. 9). Next, we measure the linewidth of the resonant peak and L–L curves in the symmetry-protected BIC lasers more systematically by varying the lattice constants (Supplementary Figs. 10 and 11). Fig. 4a shows representative measured data for three different lattice constants. Clear lasing behaviors are also observed, although the below-threshold linewidth is small at \( a = 574 \) and 578 nm due to the high Q factor, which will be discussed in Fig. 4c. The threshold is much lower in the super-BIC regime at \( a = 574 \) nm. This feature is shown more clearly in the plot of the threshold power density vs. the lattice constant (Fig. 4b). The threshold decreases significantly as the lattice
constant increases until \( a = 574 \text{ nm} \) and becomes almost constant at \( a \geq 574 \text{ nm} \). Thus, the measurement indicates that the lasing threshold is minimized by the transition to the super-BIC mode. We note that the threshold peak power of \( \sim 340 \mu\text{W} \) and threshold power density of \( \sim 14.7 \text{ W/\mu m}^2 \) at \( a = 574 \text{ nm} \) are the smallest among all the BIC lasers previously reported (Supplementary Table 1)\(^{5,13-19,29}\).

To understand the ultralow threshold of the super-BIC laser, we estimate the Q factor, \( \lambda/\Delta\lambda \), where \( \lambda \) is the peak wavelength and \( \Delta\lambda \) is the linewidth of the peak at the transparent pumping condition at \( \sim 340 \mu\text{W} \) and threshold power density of \( \sim 14.7 \text{ W/\mu m}^2 \) at \( a = 574 \text{ nm} \) are the smallest among all the BIC lasers previously reported (Supplementary Table 1)\(^{5,13-19,29}\).

Fig. 4 Threshold and Q factors of BIC lasers. a Measured L–L curves (black; left y-axis) and linewidth (blue; right y-axis) of the symmetry-protected BIC lasers at \( a = 571 \text{ nm}, 574 \text{ nm}, \) and \( 578 \text{ nm} \). The increase of the linewidth above the threshold is due to the thermal effect. b Measured threshold values divided by the pump area (\( \sim 5.4 \text{ \mu m} \times \text{ size} \)) as a function of the lattice constant. c Measured Q factors, \( \lambda/\Delta\lambda \), as a function of the lattice constant. \( \Delta\lambda \) is the linewidth estimated at the transparent pumping condition (\( \sim 340 \mu\text{W} \)). The linewidth data were taken from Supplementary Fig. 10. The orange line indicates the Q factor obtained using the resolution-limited linewidth of the spectrometer (\( \sim 0.22 \text{ nm} \)).

Our measurements show that the super-BIC mode has a higher Q factor than the other BICs before merging, as indicated by Fig. 2. The effective increase in the number of air holes confining the super-BIC mode (Fig. 3e) further enhances the Q factor. The Q factor starts to decrease again at \( a = 579 \text{ nm} \) (Fig. 4c), which indicates that the merging effect ends in this regime. This feature is more evident when the pump spot size increases (Supplementary Fig. 12). For larger pump spot sizes, the threshold values at \( a = 578 \text{ nm} \) and \( 579 \text{ nm} \) increase, whereas the super-BIC laser still shows a low threshold (Fig. 5a). In particular, the super-BIC occurs in a narrower range of lattice constants as the pump spot size increases, which is like the simulation result shown in Fig. 2e. Also, the super-BIC regime from \( a = 574 \text{ to } 577 \text{ nm} \) agrees well with the regime of calculated high radiative Q factors. Therefore, the merging point at \( a \sim 574 \text{ nm} \) is further clarified through the experiment performed with varying pump spot sizes. At \( a = 578 \text{ and } 579 \text{ nm} \), the lasing mode turns into the isolated BIC.

Furthermore, additional laser properties are investigated in the super-BIC regime (\( a = 574 \text{ nm} \)). First, we measure the polarization-resolved lasing images by placing a linear polarizer in front of the IR camera (Fig. 3b). These images exhibit an intensity minimum along the direction of the polarizer, which agrees well with the previous report\(^{21}\). Second, we estimate the spontaneous emission factor, by comparing the measured L–L curve with that obtained from the conventional rate equations (Fig. 5c). The estimated spontaneous emission factor is \( \sim 0.01 \), which is smaller than the values of ultrasmall nanolasers\(^{32}\) due to the relatively large mode volume\(^{33-35}\). Third, the interference images are measured in the spontaneous emission, amplified spontaneous emission, and lasing regions in the super-BIC laser (Fig. 5d and “Methods” section). The interference pattern is clearly observed only in the lasing region, exhibiting the calculated coherence time of \( >38 \text{ ps} \) (ref. 36). Fourth, we measure the decay times in the spontaneous emission and lasing regions of the super-BIC laser (Supplementary Fig. 13). The measured decay time of \( <138 \text{ ps} \) in lasing is fast enough for the high-speed modulation\(^{37,38}\). 
We have demonstrated a super-BIC laser based on a finite-size photonic cavity with a small footprint. We have observed a transition to the super-BIC laser from the symmetry-protected BIC and accidental BIC lasers by tuning the lattice constant. The theoretical analysis shows that the radiation loss in super-BIC regime is limited by the transparency value of the gain material, as a result of the low radiative loss in the finite photonic structure. Notably, its threshold is extremely low and is significantly reduced compared with other BIC lasers using similar gain materials. Moreover, the semiconductor active material with high optical gain supports the superior optical properties of super-BIC lasers.

For the practical implementation of such an ultralow-threshold stretchable laser structure to vary the lattice constant and find the merging point more easily. In addition, electrical pumping should be performed: an efficient current path needs to be formed to inject carriers to the whole area of the BIC cavity. We believe that our findings will pave the way to significantly reduced optical losses in active nanophotonic structures with a finite footprint and the development of an ultralow-threshold light source for photonic integrated circuits, by controlling the topological charges in the reciprocal space and engineering the radiation condition.

Methods

Numerical simulations. The photonic band diagrams and optical properties of the resonant modes are calculated in a free-standing InGaAsP membrane using a three-dimensional finite-element method (FEM) solver in COMSOL Multiphysics. Floquet periodic boundaries and perfectly matched layers are imposed in the in-plane and vertical directions, respectively, for the infinite-size structures (Fig. 1 and Supplementary Figs. 1–4). The radiation loss (γ) is calculated by collecting the out-of-plane component of the Poynting vector away from the surface of the slab (Fig. 1b). Each radiation loss is normalized by the radiative loss at the merging point more easily. In addition, electrical pumping should be performed: an efficient current path needs to be formed to inject carriers to the whole area of the BIC cavity. We believe that our findings will pave the way to significantly reduced optical losses in active nanophotonic structures with a finite footprint and the development of an ultralow-threshold light source for photonic integrated circuits, by controlling the topological charges in the reciprocal space and engineering the radiation condition.

Methods

Discussion

We have demonstrated a super-BIC laser based on a finite-size photonic cavity with a small footprint. We have observed a transition to the super-BIC laser from the symmetry-protected BIC and accidental BIC lasers by tuning the lattice constant. The theoretical analysis shows that the radiation loss in super-BIC follows the law $k^6$ depending on the lattice constant and the super-BIC keeps a high Q factor even in the finite-size cavity. Thus, the high-performance optical characteristics, including an ultralow threshold, a single lasing peak, and a high Q factor, have been measured for the super-BIC laser. These features in the super-BIC laser are distinguishable from those in the symmetry-protected BIC laser. Notably, its threshold is extremely low and is limited by the transparency value of the gain material, as a result of the low radiative loss in the finite photonic structure. Furthermore, the semiconductor active material with high optical gain well supports the superior optical properties of super-BIC. Compared with other BIC lasers using similar gain materials (Supplementary Table 1), the threshold values we have measured are lower in the super-BIC regime and similar outside the super-BIC regime.

For the practical implementation of such an ultralow-threshold super-BIC laser, it is necessary to develop a flexible and stretchable laser structure to vary the lattice constant and find the merging point more easily. In addition, electrical pumping should be performed: an efficient current path needs to be formed to inject carriers to the whole area of the BIC cavity. We believe that our findings will pave the way to significantly reduced optical losses in active nanophotonic structures with a finite footprint and the development of an ultralow-threshold light source for photonic integrated circuits, by controlling the topological charges in the reciprocal space and engineering the radiation condition.
structures (Fig. 2), perfectly matched layers are introduced in all directions including the in-plane domain of $N \times N$ size. The far-field simulation is performed using a circular-shaped outer boundary to remove artificial interference patterns (Fig. 3d).

In addition, a home-made three-dimensional finite-difference time-domain (FDTD) method is used (Fig. 2f and Supplementary Fig. 5c) to calculate the radiative Q factor and cross-check the validity of the result in Fig. 2e, because FDTD-Q factors reflect finite-size cavity in the time domain by observing the time decay of resonant modes. The convolutional perfectly matched layer is used as an absorbing boundary condition in the FDTD. The size of the mesh grid is 10 nm, and more than 1000 periods of resonance oscillations are observed in the time domain to precisely calculate resonant wavelengths and Q factors. The Poynting vector is decomposed into in-plane and vertical components in the slab structure, and the radiative Q factor is then obtained using the vertical component.

In the FEM and FDTD simulations (except for Supplementary Fig. 5), the hole diameter and slab thickness are set to 480 and 650 nm, respectively. Different structural parameters are examined in Supplementary Fig. 5. The refractive index of the InGaAsP slab is set to 3.4.

**Device fabrication.** The samples are fabricated using a 650 nm-thick InGaAsP/1 μm-thick InP/100 nm-thick InGaAs/InP substrate wafer. The InGaAsP layer includes seven 7 nm quantum wells in the middle, whose central emission wavelength varies from 1.5 μm. The InP and InGaAs layers act as sacrificial and etch stop layers, respectively. To define a periodic square-lattice structure, electron-beam lithography is performed at 30 keV on a polymethyl methacrylate (PMMA) layer coated on the wafer. The hole diameter is fixed at ~400 nm and the lattice constant varies from 560 to 580 nm. Chemically assisted ion-beam etching is performed to drill air holes in the InGaAsP layer while using the PMMA layer as an etch mask. Finally, the sacrificial InP layer is selectively wet-etched using a diluted HCl: H2O (4:1) solution at room temperature, and the remaining PMMA layer on top of the slab is removed by O2 plasma.

**Optical measurements.** A 980 nm pulsed laser diode (20% duty cycle, 1 MHz period) is used to optically pump the fabricated samples at room temperature. The light emitted from the cavities is collected by a ×100 objective lens with a numerical aperture of 0.85 (LCLPLN100XIR, Olympus) and focused on either a spectrometer with an IR array detector (SP 2300i and PyLoN, Princeton Instruments) or an InGaAs IR camera (C10633, Hamamatsu). The spot size of the pump laser is varied from ~5.4 to ~9.2 μm using additional bulk lenses. The resolution of the spectrometer is ~0.22 nm. In Fig. 3b, the wavelength is taken just above the threshold of the optical measurement. Received: 29 January 2021; Accepted: 22 June 2021; Published online: 05 July 2021

**Data availability**

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

**Code availability**

The codes used in this work are available from the corresponding authors upon request.

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Author contributions

H.-G.P. and Y.K. conceived the research and supervised the project. M.-S.H., H.-C.L., and K.-Y.J. fabricated the samples and conducted experimental studies. M.-S.H., K.-H.K., S.-H.K., and K.K. performed numerical simulations and theoretical analyses. M.-S.H., H.-C.L., K.-H.K., Y.K., and H.-G.P. wrote the manuscript based on the input from all the authors. All authors contributed to writing and editing the manuscript.

Competing interests

The authors declare no competing interests.

Additional information

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