Pushing the polariton confinement limits with low losses
using image polaritons in boron nitride

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Polaritons in two-dimensional (2D) materials provide extreme light confinement that is not possible with conventional metal plasmonics. However, such large confinement inevitably increases optical losses through various damping channels. Here we demonstrate that a hyperbolic phonon polariton (HPhP) mode in hexagonal boron nitride (hBN) can overcome the fundamental trade-off, featuring ultra-tight confinement and exceptional quality factors simultaneously. Compared to its charge-symmetric counterpart, the HPhP mode with anti-symmetric charge distributions exhibits lower optical losses while also acquiring tighter polariton confinement that is not limited by Landau damping due to the bosonic nature of HPhPs. Far-field observation of this high-momenta anti-symmetric mode becomes possible with our resonator design that can boost the coupling efficiency by launching a virtual polariton mode with image charges, while using hBN isotopically enriched in $^{10}$B with low optical losses. We experimentally observed a record-high effective index of up to 132 and quality factors as high as 501, values much higher than found in hBN with the natural distribution of boron isotopes. Our phenomenological theory suggests an important role of hyperbolic surface scattering in the damping process of HPhPs. Our image-polariton resonator scheme is universally applicable to polaritons in 2D materials and enable the extreme engineering of modal properties and light-matter interactions.

Surface phonon polaritons (SPhPs)\textsuperscript{1-3} are collective oscillations of atomic lattice vibrations coupled with electromagnetic waves, supported by polar materials such as hexagonal boron nitride (hBN)\textsuperscript{4,5}, silicon carbide (SiC)\textsuperscript{6}, and molybdenum oxide (MoO$_3$)\textsuperscript{7}. PhPs are supported within spectral regions bounded by the frequencies of the longitudinal (LO) $\omega_{\text{LO}}$ and transverse optic (TO) $\omega_{\text{TO}}$ phonon modes referred to as the Reststrahlen band, where the real part of permittivity of a polar material becomes negative. A characteristic feature of PhPs\textsuperscript{1,3,9} is their ability to confine radiation to subdiffractional length scales\textsuperscript{10,11}. In contrast to diffraction-limited electromagnetic modes in dielectric waveguides, the polariton wavelength can be continuously reduced in polariton-supporting materials. For slabs thinner than the evanescent extent of the polaritonic fields, the polariton wavelength can be strongly modified by the thickness of the polaritonic material. In this regard, two-dimensional (2D) materials offer an exciting opportunity to explore highly confined polaritons with unprecedentedly large in-plane momenta\textsuperscript{10,12-14}. Recent works on PhPs in 2D materials\textsuperscript{2-3} such as hBN have reported effective indices of a few tens - several times larger than values previously reported for plasmons in metallic nanostructures or graphene. While graphene plasmons can be more tightly confined than metal plasmons, their confinement (and hence effective index) is ultimately limited by Landau damping\textsuperscript{15-18}, through which a plasmon creates an electron-hole pair as it enters the single particle phase space. In contrast, PhPs in hBN are unimpeded by Landau damping due to the bosonic nature of both phonons and photons, suggesting that field confinement via PhPs in hBN can exceed what is possible with graphene plasmons.

Despite its distinctive physical origin, the SPhPs are phenomenologically similar to surface plasmon polaritons (SPPs)\textsuperscript{19-21}, which originate from collective oscillations of electrons. While strong confinement of polaritons can be achieved for SPP and SPhP modes near the plasma frequency $\omega_p$ and longitudinal (LO) phonon frequency $\omega_{\text{LO}}$, respectively, extreme confinement has recently been demonstrated for graphene plasmons, when placed in proximity to a metal plate with a thin dielectric spacer. The modified Coulomb interactions lead to a linear dispersion known as an acoustic plasmon\textsuperscript{22-28}. The acoustic graphene plasmons originate from coupling between the plasmons supported by a graphene layer and its mirror image inside the conducting plate with out-of-phase charge oscillations. As a result, most of the electric field is confined within the gap between graphene and the conducting plate, resulting in extreme plasmon confinement several times tighter than conventional graphene plasmons. The ultimate confinement limit of the acoustic graphene plasmon, however, is still set by Landau damping, which is given by $c/\nu_F = 300$ with $c$ and $\nu_F$ being the speed of light and the graphene Fermi velocity\textsuperscript{27}. Thus, the bosonic nature of the PhPs raises intriguing questions regarding the fundamental characteristics of such image polaritons and their ultimate limit when Landau damping is irrelevant.
Phonon-polaritons in hBN are long-lived\textsuperscript{4,29}, and can be harnessed to store electromagnetic energy in resonators\textsuperscript{10,30,31} and in propagating modes\textsuperscript{29,32,33}. Due to the opposite signs of axial and tangential permittivities in the two Reststrahlen bands of hBN, it can support hyperbolic phonon polaritons (HPhPs)\textsuperscript{10,13,34,36}. For propagating HPhPs, the figure-of-merit (FoM) can be defined as the ratio of the real part of the momentum to the imaginary part. Values up to 33 have been measured by scattering-type near-field optical microscopy\textsuperscript{13}. For HPhP resonators, quality factors of 70 and 283 have been reported for the resonators based on hBN ribbons\textsuperscript{31} and nanocones\textsuperscript{10}, respectively, which are one or two orders of magnitude larger than the quality factors of graphene plasmonic resonators\textsuperscript{28,37}. Recently, ultrapure hBN isotopically enriched in either \textsuperscript{10}B or \textsuperscript{11}B (with longer phonon lifetimes than hBN with the natural 20\% \textsuperscript{10}B and 80\% \textsuperscript{11}B isotope distribution) has been utilized to further push the FoM to beyond 40 for propagating modes\textsuperscript{29} and the quality factor to \textasciitilde300 for resonators. While these performance metrics are already impressive, their ultimate limit has yet been experimentally probed and understood. Importantly, extreme subwavelength confinement via hBN implies that free-space photons cannot excite these high-momenta HPhP modes efficiently. To construct practical polaritonic devices it is desirable to simultaneously accomplish tight confinement as well as highly efficient far-field coupling of light, while pushing the limit of quality factors unimpeded by Landau damping.

Here, we solve these challenges by using monoisotopic \textsuperscript{10}BN and our image polariton resonator with a high coupling efficiency to harness new image HPhP modes. The resonator consists of an array of metal ribbon antennas located a few nanometers away from an hBN layer to launch the polaritons. As in the acoustic plasmon case, the presence of image charges improves the confinement of the HPhPs, illustrating that such strategies can be universally applied for polaritonic modes. This high confinement was correlated with higher-order antisymmetric modes exhibiting record-high effective indices up to 132 and quality factors as high as 501, respectively. The symmetric modes are found to have relatively lower quality factors and smaller effective indices compared to the anti-symmetric counterpart, but exhibit higher coupling efficiency allowing for the extinction of nearly 90\% of incident light. Based on the strong dependence of the quality factor upon the dielectric gap thickness, we suggest surface scattering effects as one of the primary damping pathways and develop the phenomenological theory for describing the loss mechanism of the HPhPs. The high quality factors achieved for such large effective indices are unprecedented compared with other electromagnetic modes such as SPPs in metals and graphene, where the increase in the confinement is compensated by increases in the optical loss.

**Concept of image polaritons**

In the presence of a mirror plane a few nanometers away from a polariton-supporting material, a polariton mode is mirrored, but with an inverted charge distribution (Fig. 1a). This strategy has been utilized to enhance the confinement of SPPs in noble metals\textsuperscript{38,39}. In the case of hBN, HPhP modes supported in a slab of thickness \(t\) are quantized in the out-of-plane direction (\(y\)-axis) with the out-of-plane wavevector of \(k_y = \pi / t\), where \(l\) denotes the mode number. For very thin films, coupling efficiencies to modes with large \(l\) become negligible due to the large momentum mismatch and only a couple of modes with small \(l\) are reasonably stimulated. Let us consider the case where only two modes with \(l = 1\) and 2 are practically accessible. The 1st mode with the polaritonic fields along the in-plane direction oscillating in-phase is named the ‘symmetric mode’ (Fig. 1b), while the 2nd mode with out-of-phase charge oscillations is ‘anti-symmetric mode’ (Fig. 1c) (see Fig. S1). Due to the out-of-phase charge oscillation, most of the electric fields for the anti-symmetric mode are confined within the hBN slab, leading to stronger polariton confinement than for the symmetric mode. Due to the opposing charge distribution between the original and mirror polaritons, the electric fields are tightly confined within a dielectric gap between the hBN layer and conducting plane. The presence of a conducting plane can better confine the in-plane HPhPs of different
modal nature with the confinement increasing with decreasing dielectric gap size (Fig. 1d) (see supplementary note 1). Also, the in-plane momentum increases as the frequency of the incident wave approaches \( \omega_{\text{LO}} \) in hBN. While in graphene the confinement is limited to \( c/\pi \approx 300 \) due to Landau damping, and the maximum quality factor by electron-phonon scattering\(^{37,40}\), the confinement in hBN image PhPs is limited only by the intrinsic material loss. These image phonon-polaritons exhibit enhanced FoMs as the gap size scales down to the nanometer regime, and can also surpass that of its SPP analog in graphene (Fig. 1e). Here, we realize a highly efficient polaritonic nanoresonator with isotopically enriched hBN to uncover the experimental upper bounds of these image phonon-polaritons, akin to the acoustic polaritons in graphene-based devices.

**Gap dependence and dispersion**

Far-field observation of image polaritons becomes possible with our resonator design illustrated in Fig. 2a. In contrast to previous resonator designs for HPhPs, our approach utilizes an unpatterned hBN flake, minimizing the scattering of the polaritons at the patterned edges of resonators and eliminating additional damages induced through traditional lithographic patterning. Instead, the continuous conducting plate that was previously used for mirroring the polaritons as in Figs. 1b and 1c is patterned into an array of optical nanoribbon antennas to launch polaritons. The substrate that includes the metal ribbons, optical spacer, and reflector is fabricated using a template-stripping process\(^{28,41}\), which produces an ultraflat top surface even in the presence of the undulating topography of the metal ribbons. The optical spacer and the reflector further boost the coupling efficiency by recoupling the transmitted waves back to the image polaritons when the thickness of the optical spacer is designed to satisfy the quarter-wavelength condition\(^{28}\). The optical microscopic image in Fig. 2b shows the top view of the resonator with the hBN layer on top of metal ribbon patterns. The scanning electron microscope image in Fig. 2c shows that the top surface of the metal ribbons (\( p = 150 \, \text{nm} \)) is smooth and free from defects over a large area.

The gap dependence of the image polaritons can be studied using Fourier transform infrared (FTIR) spectroscopy. The reflection spectra measured from devices with three different gap sizes of 3, 8, and 20 nm given the metal ribbon periodicity (\( p \)) of 150 nm are shown in Fig. 2d. The large resonance dips at around 1500 cm\(^{-1}\) originate from the symmetric mode. Particularly for \( g = 8 \, \text{nm} \), the resonant absorption is nearly 90\%, showing the high coupling efficiency of the image polariton resonator enabled by the two-stage coupling mechanism and the photon-recycling scheme\(^{28}\). As the gap size decreases, the resonance shifts to lower frequencies at a given in-plane momentum. On the other hand, the smaller amplitude resonances at around 1400 cm\(^{-1}\) corresponds to the anti-symmetric modes, which are less sensitive to the gap size. The spatial distribution of electric fields for the symmetric mode in Fig. 2e shows characteristic hyperbolic rays inside the hBN layer and highly confined fields within the gap region. For the anti-symmetric mode (Fig. 2f), the field intensities inside the hBN layer and the gap are of similar magnitudes. The larger intensities inside the hBN layer are associated with the out-of-phase oscillations of polaritonic fields on the top and bottom surface of the hBN layer. Such spatial distributions also explain why the anti-symmetric mode is less sensitive to the gap size. To investigate the gap-dependence of the resonant absorption, the resonances with similar in-plane momenta as well as resonance frequencies are compared. Instead of resonant absorption itself, the magnitudes of the Lorentzian fits for the resonances are plotted for fair comparison. For the symmetric case, the resonances at around 1522 cm\(^{-1}\) are considered (Fig. 2g), while for the anti-symmetric case, the resonances at around 1416 cm\(^{-1}\) are selected. Due to the different modal natures of the two modes, the resonant absorption scales differently with the gap size (Fig. 2h). For the anti-symmetric mode, placing the metal ribbon array closer to the hBN layer increases resonant absorption due to the rapid decay of the scattered fields and the resultant sharp contrast in the intensity on the top and bottom surfaces of the hBN layer. On the other hand, the reduced separation between hBN and the metal
ribs will decrease the resonant absorption of the symmetric mode, the excitation of which requires uniform field intensities along the out-of-plane direction. These qualitative trends are well captured by our simulation, albeit with a constant scaling factor (<1) to the simulation results (see Fig. 2h caption). The decreased absorption is understandable due to known defects such as grain boundaries, and non-uniform thickness of hBN flakes as well as imperfections in metal nanostructures. The defects in hBN impact the anti-symmetric mode more severely due to the confined fields within hBN. From spectra measured from devices with different \( p \), the in-plane momenta are extracted using the approximate equation \( k = 2\pi n/p \), where \( m \) is the order of resonance and the effective indices are obtained by \( k/k_0 \), with \( k_0 \) being the free-space wave-vector (Fig. 2i) (for all the measured spectra, see Figs. S2 and S3). The experimental results agree well with the theoretical dispersions and show two distinct branches, namely the symmetric and anti-symmetric modes. For both modes, the in-plane momentum increases as the gap size decreases, though the anti-symmetric mode is less sensitive. Our strategy of using image polaritons improves the polariton confinement, producing record-high effective indices of 132 and 85 for the anti-symmetric and symmetric modes, respectively, the former of which is among the highest values reported for all polariton modes supported in polar dielectrics such as hBN, MoO\(_3\), and SiC.

**Naturally abundant vs. isotopically enriched hBN**

The utilization of hBN enriched with \( ^{10}B \) isotopes made possible the observation of modes with inherently low spectral weights. The theoretical polariton dispersions for naturally abundant hBN as well as h\(^{10}\)BN are provided in Fig. 3a. These dispersions coincide after the offset between the Reststrahlen bands for the two forms of hBN is compensated by a spectral shift for the naturally abundant hBN by 34 cm\(^{-1}\) towards higher frequencies. The far-field spectra measured from the resonators using the natural hBN and h\(^{10}\)BN with nominally identical geometrical dimensions are shown in Fig. 3b. After compensating for the offset, the resonances for the symmetric mode measured from the naturally abundant hBN devices with different gap sizes are well aligned to those of the h\(^{10}\)BN structures, consistent with Fig. 3a. The quality factors of the resonances for the symmetric mode measured from the resonators with the naturally abundant hBN are 60, 75, and 155 for \( g=3, 8, \) and 20 nm, respectively. For the h\(^{10}\)BN case, much larger quality factors of 81, 137, and 209 were measured. For all the gap sizes, the resonant absorption of the symmetric mode for the h\(^{10}\)BN case is much higher than the resonant absorption with normal hBN due to the larger internal resonance enhancement given as \( 1 - t_{12}t_{21} \exp(-2\pi/\text{(FoM)}) \), where \( t_{12} \) and \( t_{21} \) represents the modal transmission coefficients of the HPhPs from the metal-free region to the metal-coupled region and vice versa, respectively (see supplementary note 2). Due to the larger \( t_{12} \) and \( t_{21} \) for the anti-symmetric mode (see Fig. S4), the internal resonance enhancement factor for the h\(^{10}\)BN can be an order of magnitude larger than naturally abundant hBN. The larger internal resonance enhancement makes resonances from the anti-symmetric mode in h\(^{10}\)BN clearly observable, which are difficult to resolve for naturally abundant hBN due to the small spectral weight of the anti-symmetric mode.

**Loss mechanism**

Due to the zigzag propagation of HPhPs in a hBN slab, the surface scattering is expected to play an important role on the damping of the HPhPs. The rate of the surface scattering increases with the hyperbolic ray propagation angle \( \theta_{ph} \) of the HPhPs as shown in Fig. 4a, which is given as \( \arctan \left( \sqrt{\frac{\varepsilon_f(\omega)}{\varepsilon_t(\omega)}} \right)^{10} \). Based on this concept, we have developed an analytical model for the scattering rate \( \gamma_{\text{total}} \) of the image polariton based on the loss equations for conventional Fabry-Perot resonators\(^1\), which is given as

\[
\gamma_{\text{total}} = \gamma_e + \gamma_p + \gamma_s = -\frac{\ln (t_{12}t_{21})}{\tau_p} + 2Im(k)v_g + \frac{v_g |\tan(\theta_{ph})|}{t} \frac{A}{1 + (\frac{A}{g_c})}.
\]  

(1)
where $\gamma_e$, $\gamma_p$, and $\gamma_s$ represent the scattering rates for three major damping pathways considered in this model; the scatterings at the interface between the resonator units, propagation loss, and hyperbolic surface scattering (see supplementary note 2). $T_{12}$ and $T_{21}$ represent the modal transmittances of the image polariton through the resonator units given as $|t_{12}|^2$ and $|t_{21}|^2$, respectively. $\tau_p$ and $v_g$ denote the dwelling time within a resonator unit and the group velocity of the phonon polariton. $A$ is a dimensionless fitting parameter that represents the intensity of the scattering while $g_c$ is a fitting parameter for a critical gap size. The surface scattering model is inversely proportional to the gap size based on the fact that the hyperbolic rays become more evident as the gap size decreases. The best fit with the experimental results is obtained when $A = 0.118$ and $g_c = 21.03$ nm. The total scattering rates estimated from our analytical model for $g = 3$ nm agree well with the experimental results as shown in Fig. 4b (see the anti-symmetric case in Fig. S5). Due to the long lifetime of the HPHPs in $^{10}$BN, the contribution from the propagation loss is insignificant. The surface scattering is most severe at frequencies near the middle of the Reststrahlen band. The decrease in the contribution from the surface scattering near $\omega_{TO}$ is attributed to small $\theta_{ph}$ while in the spectral vicinity of $\omega_{LO}$, this is due to small $v_g$. Edge scattering plays an important role at around $\omega_{TO}$, where $T_{12}$ and $T_{21}$ become small due to the smaller effective index of the image polariton.

Our model describes the trend of the quality factor as a function of the device gap and other physical parameters with good accuracy. The quality factors extracted from the resonances for the symmetric mode agree well with the theoretical results calculated using $\omega/\gamma_{total}$ (Fig. 4c), supporting our conjecture that the total scattering increases with decreasing the gap size (see Fig. S6 for the extraction method of the quality factor). The increase in the quality factor at higher frequencies is also well explained with the theory that associate the lower damping with the decrease in $v_g$. The theoretical model predicts that the quality factors for the anti-symmetric mode are larger than the quality factors for the symmetric mode due to the smaller group velocity and larger modal transmittances. Indeed, the quality factors measured for the anti-symmetric mode are much larger than the quality factors for the symmetric mode (Fig. 4d). The fit with the experimental results is best when $A = 0.108$ and $g_c = \infty$, which shows that the anti-symmetric mode is less sensitive to the gap size. In general, the decrease in a quality factor with increasing frequency is well captured by the theory. However, the experimental results show noticeable departure from the theoretical results, which can be attributed to the sensitive nature of the anti-symmetric mode to defects and material composition of hBN. Also, around 1500 cm$^{-1}$ where the polariton wavelength of the anti-symmetric mode becomes extremely small, the edge or surface scatterings seem to be more suppressed than anticipated from the theory. The numerical results also predict such high quality factors around 1500 cm$^{-1}$ (Fig. S7). Experimentally, we observed the second-order resonance of the anti-symmetric mode located at a frequency of 1507 cm$^{-1}$ with the effective index of 132 and the quality factor of 461. The highest quality factor was 501 near $\omega_{TO}$, which is the among the highest value reported for any polaritonic resonators.

Conclusions

We experimentally observe strong resonances from image phonon-polaritons in $^{10}$B isotopically enriched hBN and identified the symmetric and anti-symmetric modes that simultaneously achieve ultra-tight field confinement and high quality factors. Far-field observations of such ultra-confined modes become possible with our image polariton resonator that utilizes a pristine unpatterned hBN flake coupled with an array of ultraflat metallic ribbons, separated by a thin dielectric layer. The image polariton resonator features the high coupling efficiency and the precise control over the gap size with nanometer accuracy, which proved essential to observe the anti-symmetric mode. Most importantly, the use of unpatterned $^{10}$BN flakes with low optical losses allow us to clearly observe the anti-symmetric mode, which due to low spectral weights, has extremely high quality factors of up to 501, which was difficult to detect with natural hBN. The symmetric modes are found to have lower quality factors of up to 262 and effective indices of up to 85, yet
exhibit higher coupling efficiencies to the incident field, approaching 90%. From the measured quality factors and the effective indices, we estimated that the Purcell factor could reach $1.13 \times 10^7$, which is an order of magnitude larger than estimated for graphene$^{43}$ and metal$^{44}$ plasmons at mid-infrared frequencies. The theoretical model developed here provides a clear understanding of the loss mechanisms and a strategy to design optimal polaritonic nanoresonators. Our far-field image polariton platform and new understanding of the loss mechanism enable researchers to harness ultralow-loss and ultra-compressed cavity polariton modes inaccessible via graphene or conventional metal plasmons. This capability, in turn, will enable a series of new fundamental studies that requires ultrastrong-light matter interactions, as well as nonlinear$^{45,46}$ and nonlocal$^{27}$ effects. Also, the ability of our platform to efficiently couple to far-field light will be essential for the development of high-contrast polaritonic sensors, optoelectronic devices, and emitters. While plasmons in graphene or noble metals have been considered the most promising route to extreme subwavelength confinement, our findings show that HPhPs in hBN can outperform plasmons for a broad variety of applications.

Author contributions
I.-H.L. and S.-H.O. conceived the idea. J.H.E. and S.L. synthesized h$^{10}$BN crystals. I.-H.L., M. H., X. Z., and Y. L. performed device fabrication. I.-H.L. and M. H. characterized devices. I.-H.L., T.L., P.A., K.W., J.D.C. and S.-H.O. performed theoretical analysis. All authors analysed the data and wrote the paper together.

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Competing interests. The authors declare no competing interests.
Methods

**Hexagonal boron nitride crystal growth.** The monoisotopic hBN crystals were grown as described elsewhere. In brief, 4.0 weight percent of boron which was 99.22 % the B-10 isotope was mixed with equal masses of nickel and chromium, then heated to 1550 °C to form a liquid solution. This was performed at atmospheric pressure under a nitrogen plus hydrogen gas mixture. The solution was slowly cooled which reduced the hBN solubility, causing crystals to form on the surface. These were removed from the boule with tape, after the boule reached room temperature.

**Device fabrication.** A (100) silicon wafer was heated on a hot plate at 180°C for 5 min and treated using a standard oxygen plasma recipe for 2 min (Advanced Vacuum, Vision 320). A 60-nm-thick gold film was deposited as a sacrificial layer using an electron-beam evaporator (CHA Industries, SEC 600). Then, an alumina layer was deposited using an atomic layer deposition (ALD) system (Cambridge Nano Tech Inc., Savannah). On top of the alumina film, a 30-nm-thick gold film was deposited using plasma sputtering (AJA, ATC 2200). The metal film was patterned by using an electron-beam lithography system (Vistec, EBPG5000+) followed by an Ar⁺ ion milling process (Intlvac, Nanoquest) at a beam current of 70 mA and an accelerating voltage of 24 V for 4 min. The remaining PMMA resist was removed by oxygen plasma treatment. The resultant metal ribbons were covered by alumina, silicon oxide, amorphous silicon, titanium, and gold layer using ALD, plasma-enhanced chemical vapor deposition (Plasma-Therm, PECVD), and sputterer, respectively. The deposited film thicknesses were 20, 1, 530, 1, and 60 nm, respectively. The fabricated multilayer stack was peeled off from the silicon wafer and transferred to a glass substrate by using a photocurable epoxy (Norland, NOA 61) to bond the two substrates. In the transferred device on the glass substrate, the order of the whole stack is inverted, meaning that the top surface of the template-stripped device was previously the bottom interface between the silicon wafer and sacrificial gold layer. The sacrificial layer was simply removed using a gold etchant (Sigma Aldrich) with high etching selectivity to the underlying alumina film. Lastly, the mechanically exfoliated hBN flakes were the atomically clean exfoliated hBN flakes with desirable thickness were identified with atomic-force microscope and subsequently transferred onto the template-stripped substrate via a standard dry-transfer method.

**Device characterization.** The reflection spectra were measured using an optical microscope coupled to a Fourier transform infrared spectrometer (Vertex; Bruker) using a 15× Cassegrain objective lens. The spectral resolution was 2 cm⁻¹. The reflectance spectra were referenced to an aluminum mirror.

**Numerical simulation.** The spatial distributions of the electric fields were calculated using COMSOL Multiphysics (COMSOL Inc.).

**Data availability**

The data that support the plots within this paper and other findings of this study are available from the corresponding author upon reasonable request.

**Additional Information**

Supplementary information is available in the online version of the paper. Reprints and permission information is available online at www.nature.com/reprints. Correspondence and requests for materials should be addressed to S.-H.O.
Fig. 1. Concept of image polaritons. a, Schematic illustrations of image polaritons. Charge distributions for b, symmetric mode and c, anti-symmetric mode. In the presence of a mirror, image polaritons with out-of-phase charge oscillations are launched and confine radiation to the gap between the hBN layer and the mirror. d, The normalized in-plane momentum and e, figure of merit as a function of the gap size for the symmetric and anti-symmetric polariton modes, and graphene plasmons at five different frequencies of 1522, 1545, 1569, 1592, 1616, and 1639 cm\(^{-1}\). The arrows represent the direction of increasing frequency. For graphene, the doping level and damping rate are 0.4 eV and 0.00378 fs\(^{-1}\), respectively.
**Fig. 2. Gap dependence and dispersion.** 

**a,** Schematic illustration of the image polariton resonator. \( w, s, g, \) and \( t \) denote the width of a metal ribbon, the spacing between neighboring metal ribbons, the gap size between the hBN and the conducting plate, and the thickness of the hBN. \( d \) denotes the thickness of an optical spacer, which is determined by the quarter wavelength condition to enhance the coupling efficiency from far-field light to image polaritons. 

**b,** Optical microscope image of a hBN flake transferred on the template-stripped substrate. 

**c,** Scanning electron microscopy image of metal ribbons with \( p = 150 \) nm. 

**d,** The reflection spectra for \( g = 3, 8, \) and \( 20 \) nm given the periodicity of 150 nm. The spatial distribution of the electric fields for **e,** the symmetric mode and **f,** anti-symmetric mode. The white arrows represent Poynting vectors. 

**g,** Reflection spectra measured from the devices whose resonance wavelengths for the symmetric mode are almost aligned at the wavenumber of 1520 cm\(^{-1}\). 

**h,** The gap dependence of the magnitude of a resonance as a function of the gap sizes for the symmetric and anti-symmetric mode at around the frequency of 1520 and 1416 cm\(^{-1}\), respectively. The solid lines denote the numerical results with constant scaling factors of 0.75 and 0.2 for the symmetric and anti-symmetric case. 

**i,** The in-plane momenta extracted from the measured spectra for different gap sizes of \( g = 3 \) (red circle), 8 (blue triangle), and 20 nm (black square) together with the analytical dispersions. For each gap size, the dispersion is calculated for \( t = 20 \) and 30 nm (the shaded area between two solid lines) to account for hBN thickness \((t)\) variation in the fabricated samples.
Fig. 3. Naturally abundant vs $^{10}$B-enriched hBN. a, The dispersion for the naturally abundant and isotopically enriched hBN for the gap size of 3 nm. b, The measured reflection spectra for the two hBN for different gap sizes of 3, 8, and 20 nm given the periodicity of 150 nm. In all cases, the spectra from naturally abundant hBN cases are shifted toward higher frequencies by 34 cm$^{-1}$ to align the Reststrahlen bands for better visibility. The orange shaded region shows the presence of the resonances from the anti-symmetric modes, which can be observed only in the h$^{10}$BN cases.
Fig. 4. Loss mechanism. a, Schematic illustration for the surface scatterings mediated by the hyperbolic rays inside hBN. b, The loss contributions from three different damping mechanisms calculated using the analytical model in Eq. (1) for $g = 3$ nm. The red circles are the total scattering rates measured from the samples with different periodicities. The experimental (symbols) and analytical results (solid lines) for the quality factors obtained from the resonances for c, the symmetric mode and d, anti-symmetric mode.
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Supplementary Information for

Pushing the Polariton Confinement Limits with Low Losses using Image Polaritons in Boron Nitride

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Supplementary note 1. Image polariton dispersion

The magnetizing fields, $H(x, y, z)$, in medium I ($y < 0$), II ($0 < y < t$), and III ($0 < y < t$) (illustrated in Fig. S1a) can be represented as

\[
H_I(x, y, z) = \hat{e}_z (H_1 e^{iky,1y} + H_t e^{-iky,1y}) e^{ikx},
\]
\[
H_{II}(x, y, z) = \hat{e}_z (H_{2p} e^{iky,1y} + H_{2n} e^{-iky,1y}) e^{ikx},
\]
\[
H_{III}(x, y, z) = \hat{e}_z (H_{3p} e^{iky,1y} + H_{3n} e^{-iky,1y}) e^{ikx},
\]

where $k$ and $k_{y,i}$ are the wavenumber components in the $x$ and $y$ direction with the subscript $i \in \{I, II, III\}$ representing the medium of interest. For the medium I and III, $k_{y,i} = \sqrt{\varepsilon_1 (k_0^2 - \alpha^2)}$ while for the medium II, $k_{y,II} = \sqrt{\varepsilon_\parallel k_0^2 - (\varepsilon_\parallel / \varepsilon_\perp) k^2}$ with $\varepsilon_i$ and $k_0$ being the complex permittivity of the medium and the free-space wavenumber, respectively. The electric fields are given as

\[
E_I(x, y, z) = -\frac{1}{i\omega\varepsilon_1} \left[ \hat{e}_z ikH_1 + \hat{e}_x (ik_{y,1}H_1 e^{iky,1y} - ik_{y,1}H_t e^{-iky,1y}) \right] e^{ikx},
\]
\[
E_{II}(x, y, z) = -\frac{1}{i\omega\varepsilon_\parallel} \left[ \hat{e}_z ikH_{II} + \hat{e}_x (ik_{y,II}H_{2p} e^{iky,IIy} - ik_{y,II}H_t e^{-iky,IIy}) \right] e^{ikx},
\]
\[
E_{III}(x, y, z) = -\frac{1}{i\omega\varepsilon_{III}} \left[ \hat{e}_z ikH_{III} + \hat{e}_x (ik_{y,III}H_{3p} e^{iky,IIIy} - ik_{y,III}H_{3n} e^{-iky,IIIy}) \right] e^{ikx}.
\]

From the continuities of the tangential components of the fields at the three interfaces ($y = 0, t$, and $g$), we have

\[
\begin{align*}
H_I + H_t &= H_{2p} + H_{2n}, \\
H_I - H_t &= \alpha_1 (H_{2p} - H_{2n}), \\
H_{2p} e^{iky,IIy} + H_{2n} e^{-iky,IIy} &= H_{3p} e^{iky,IIIy} + H_{3n} e^{-iky,IIIy}, \\
H_{2p} e^{iky,IIy} - H_{2n} e^{-iky,IIy} &= \alpha_2 (H_{3p} e^{iky,IIIy} + H_{3n} e^{-iky,IIIy}), \\
H_{3p} e^{iky,IIIy} &= H_{3n} e^{-iky,IIIy}. 
\end{align*}
\]
From the boundary conditions, the Fresnel reflection coefficient is given as

$$r_k = \frac{H_k}{H_i} = \frac{[(1-\alpha_1)e^{-ik_y\Pi t} M + (1+\alpha_1)e^{ik_y\Pi t} N]}{[(1+\alpha_1)e^{-ik_y\Pi t} M + (1-\alpha_1)e^{ik_y\Pi t} N]}.$$  \hspace{1cm} (4)

$$M = (1 - \alpha_2)e^{i2k_y\Pi t}(g-t) + (1 + \alpha_2),$$

$$N = (1 + \alpha_2)e^{i2k_y\Pi t}(g+t) + (1 - \alpha_2).$$

The dispersion of the image polariton is obtained when the denominator vanishes as follows.

$$[(1 + \alpha_1)e^{-ik_y\Pi t} M + (1 - \alpha_1)e^{ik_y\Pi t} N] = 0.$$ \hspace{1cm} (5)

Figure S1b shows the imaginary parts of $r_k$ as a function of frequency and momentum. Similar to the case without a conducting plate, the dispersion consists of distinct branches that originate from the quantization of the out-of-plane wavevector. The spatial distributions of the $x$ components of the electric fields for the first (Fig. S1c) and second (Fig. S1d) branch show the symmetric and anti-symmetric nature of the two modes, respectively.

**Fig. S1. Dispersion and mode symmetry.** **a**, A schematic illustration for the geometry of interest. **b**, The imaginary parts of the Fresnel reflection coefficient $r_k$ as a function of frequency and momentum. The spatial distributions of the $x$-components of the electric fields for the mode **c**, on the first branch and **d**, on the second branch.
Supplementary note 2. Analytical model for resonator loss

Compared to a conventional Fabry-Perot resonator where light is mostly reflected at the ends of a resonator unit, the polariton is mostly transmitted from one resonator unit to its neighbors at the ends of a resonator unit in our case. Thus, the polaritons transmitted from a neighboring resonator unit constructively interfere with the polariton coupled from incident light to the resonator unit at destination. In this case, the ratio of the resonantly enhanced field $E_{\text{res}}$ to the initial field without interference $E_0$, is given as

$$\frac{E_{\text{res}}}{E_0} = \frac{1}{1-t_{12}t_{21}\exp[i(k_2w+k_1s)]},$$  \hspace{1cm} (6)

where $t_{12}$ and $t_{21}$ are modal transmission coefficients from the metal-free region to the metal-coupled region and vice versa, respectively. $k_2$ and $k_1$ are the momenta for the image and normal polariton supported by the metal-coupled and metal-free region, respectively. The square of the magnitude of Eq. (6) gives an internal resonance enhancement factor when the resonance condition given as $k_2w + k_1s = 2m\pi$ is met with $m$ being the order of resonance. The total damping rate $\gamma_{\text{total}}$ including the scatterings at the edges of a resonator unit $\gamma_e$ and the propagation loss $\gamma_p$ can be modelled from the relation of

$$\gamma_{\text{total}} = \gamma_e + \gamma_p \approx -\frac{\ln(T_{12}T_{21})}{t_p} + 2v_g Im(k_z).$$ \hspace{1cm} (7)

where $T_{12}$ and $T_{21}$ are modal transmittances given as $T_{12} = |t_{12}|^2$ and $T_{21} = |t_{12}|^2$. $v_g$ is the group velocity of the polariton and $t_p$ is the dwelling time of the polariton within a resonator unit defined as $p/v_g$ with $p$ being the periodicity of the metal ribbon array. Due to the finite thickness of hBN, the polariton also undergoes scattering at the surface of hBN, which is expected to become intense for small thicknesses ($t$) and large hyperbolic angles ($\theta$) with respect to the in-plane direction. Also, the smaller distance between hBN and the metal ribbon array increases the intensities of
hyperbolic rays, which leads to more intense surface scattering. These observations lead to the addition of the surface scattering component $\gamma_s$ to the total scattering rate as follows.

$$\gamma_{total} = \gamma_e + \gamma_p + \gamma_s \approx -\frac{\ln(T_{12}T_{21})}{\tau_p} + 2v_g \text{Im}(k_2) + \frac{v_g |\tan(\theta)|}{t} \frac{A}{1 + (g/g_c)}. \quad (8)$$

Fig. S2. Measured spectra. The spectra measured for a, $g = 3$ nm with the periodicities ($p$) of 100, 120, and 150 nm, b, $g = 8$ nm with $p$ of 100, 120, and 150 nm, and c, $g = 20$ nm with $p$ of 120, 150, and 240 nm.
Fig. S3. Simulated spectra. The spectra calculated for \( a, g = 3 \text{ nm} \) with the periodicities \( (p) \) of 100, 120, and 150 nm, \( b, g = 8 \text{ nm} \) with \( p \) of 100, 120, and 150 nm, and \( c, g = 20 \text{ nm} \) with \( p \) of 120, 150, and 240 nm.

Fig. S4. Modal transmittances. \( a \), A schematic illustration of the geometry considered to calculate the modal transmittance of the polariton at the interface between two conducting plates. The modal transmittances as a function of a frequency for \( b \), the symmetric mode and \( c \), the anti-symmetric mode for different gap sizes of 3, 8, and 20 nm.
Fig. S5. Damping channels for anti-symmetric mode. The loss contributions from the scatterings at the edges of resonator units, the propagation loss, and the surface scatterings calculated from the analytical model (solid lines) together with the experimental results (red circles) for $g = 3$ nm. Compared to the symmetric case, the edge and surface scattering rates are much lower due to the tighter confinement of the anti-symmetric mode.
Fig. S6. Quality factor extractions. a, The measured reflection spectrum for the resonator with $g = 3$ nm and $p = 120$ nm. Three resonances are observed, which corresponds to the first and second order anti-symmetric resonance (A11 and A12), and the first order symmetric resonance (S11). The broad peak at 1370 cm$^{-1}$ arises from coupling between the transverse optical phonon and the cavity mode in the optical spacer, which is denoted as ‘TO’ for convenience. b, The Lorentzian fitting of b, S11 c, A11, and d, A12. In a, the resonance is first isolated from the spectrum and then fitted with a Lorentzian function using the Levenberg-Marquardt method$^3$ implemented in the commercial software (Origin; OriginLab Co.). In c where A11 is spectrally overlapped with the TO resonance, the overlapped spectra region is fitted with two Lorentzian curves using the same method. When the resonance is on a linear tilted slope baseline as in d, the isolated spectrum is fitted with a Lorentzian curve after baseline correction.

\[ y = \frac{B x}{1 + (x - x_0)^2 / w^2} \]
Fig. S7. Numerically calculated quality factors. The quality factors extracted from numerically calculated reflection spectra for a, the symmetric modes and b, the anti-symmetric mode for different gap sizes of 3, 8, and 20 nm. The solid and half-filled symbols represent the 1st and 2nd order resonances, respectively. In general, the quality factors for the 1st and 2nd resonances are of similar magnitudes.

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