Self-coupled photonic resonators made of exciton materials have recently provoked great interest in the context of light–matter interactions due to their ability to produce large normal mode splittings. In order to obtain giant Rabi energy, it is rather necessary to ensure large electromagnetic fields within exciton materials. Here, using two independent numerical algorithms, namely, the finite-element method and the rigorous coupled wave analysis, we demonstrate that, even with a moderate oscillation strength, giant Rabi splittings in excess of 250 meV can be achieved in subwavelength perovskite-based photonic crystals. This can be attributed to the fact that quasi-guided resonance modes supported by photonic systems are strongly confined inside the exciton material, highly conducing to increasing the volume of light–matter interaction. We reveal how the oscillator strength of excitons and the thickness of perovskite photonic crystals influence photon–exciton couplings. Moreover, the perovskite nanostructures investigated allow us to engineer polaritonic dispersions with linear or slow-light characters. These findings show that perovskite-based photonic crystals could be an appealing and promising platform in realizing polaritonic devices.

1. Introduction

Strong coupling between excitons and optical cavities occurs when the energy exchange rate between the two constituent states exceeds the decoherent process of individual state, which leads to the formation of new eigenstates called polaritons. Such states enable realization of some fundamental physical effects such as Bose–Einstein condensation [1] and superfluidity [2], and various applications with all-optical polaritonic devices for lasing [3], photovoltaics [4], sensing [5–7], and integrated photonics [8]. Strong exciton–photon interaction is spectroscopically manifested as an avoided crossing in the energy-momentum space (referred to as the vacuum Rabi splitting), and the interaction strength is typically represented by the Rabi frequency. Of late, demonstration of Rabi splitting with the use of many exciton platforms, such as transition metal dichalcogenides (TMDs), perovskites, nanotubes, molecules, or quantum dots, has been widely reported for a variety of optical cavities constituting photonic or plasmonic and hybrid plasmonic–photonic medium [9–15]. In above-mentioned studies, strong coupling can be established when well-defined resonant optical modes supported by external optical components are resonant with excitons.

However, recent reports have demonstrated that semiconducting exciton materials themselves can be viewed as independent optical resonators due to their high refractive indices. In this situation, the optical modes supported by the material itself can self-hybridize with the exciton modes within the same material. For instance, bulk TMD flakes can serve as a natural low-quality-factor dielectric cavity, in which the Fabry–Pérot mode strongly coupled to TMD excitons is observed in the visible spectral range [16, 17]. Furthermore, nanostructuring of TMDs has received a significant amount of attention in recent years,
because it provides more degrees of freedom in tailoring the optical properties of exciton dielectrics. By far, various TMD-based nanostructures have been demonstrated, including ultrathin photonic crystals [18] and nanostructures depending on the light–matter interaction [19–21]. For example, the strong anapole–exciton coupling was revealed in TMD-based nanoresonators [19], showing a Rabi splitting of up to 190 meV. The exciton interaction with guided-mode resonances in ultrathin TMD photonic crystals yielded a Rabi splitting of about 100 meV [21]. In terms of self-coupled TMD photonic crystals, due to their refractive index higher than 4, the fundamental waveguide modes are accessible in the excitonic regime only when the thickness of the exciton material is reduced to dimensions ranging from a few to tens of nanometers. Although this will lead to an increase of the evanescent leakage, the volume of light–matter interaction is substantially decreased, thus undesirable to observation of large normal mode splitting.

Inorganic–organic lead halide perovskites have recently been another class of excellent semiconducting excitonic materials, which are capable of addressing well these issues TMDs encounter. First of all, various TMD-based nanostructures have been demonstrated, including ultrathin photonic crystals [18] and (figure 1(a)). Media 1, and 3 form the cladding layers, and the medium 2 forms the core layer having a thickness of

$$t =$$

. The waveguiding film supports a discrete number of waveguide modes that always lie below the air light-line. For transverse-electric (TE) polarized waveguide modes, they comply with the characteristic equation

$$\gamma_i t = \tan^{-1}(\gamma_1/\gamma_2) + \tan^{-1}(\gamma_3/\gamma_2) + m\pi,$$

with non-negative integer $m$ denoting the order of waveguide mode. The parameters $\gamma_1$, $\gamma_2$, and $\gamma_3$ are

$$\gamma_1 = k_0(\beta^2 - n_i^2)^{1/2}, \quad \gamma_2 = k_0(n_i^2 - \beta^2)^{1/2},$$

and

$$\gamma_3 = k_0(\beta^2 - n_i^2)^{1/2},$$

where $\beta$ is the wavevector of the guided modes and $k_0$ is the wavevector of photons in free space. Here, we use the phenethylammonium lead iodide (PEPI) perovskite, with chemical formula $(C_6H_5C_2H_4NH_3)_2PbI_4$, as the core of the waveguide. Its dielectric function can be artificially represented as [25]

$$\varepsilon_{\text{PEPI}}(E) = n_{\infty}^2 + \frac{A}{E^2 - E^2 + i\Gamma_X E}, \quad (1)$$

where $n_{\infty}$ is the refractive index of the passive medium, $A = 0.4 \text{ eV}^2$ is the oscillator strength of PEPI exciton, $E_X = 2.4 \text{ eV}$ is the energy of exciton resonance, and $\Gamma_X = 30 \text{ meV}$ is its linewidth. In the following, when taking the Lorentz term out, the structure is passive, and otherwise is active. The slab waveguide sits in an asymmetric dielectric environment with $n_1 = 1$ and $n_3 = 1.46$ respectively corresponding to air and silica.

We calculate in figure 1(b) the $E$–$\beta$ dispersion curves of TE-guided mode in the PEPI slab waveguide. The passive structure with $t = 300$ nm supports three guided modes TE0, TE1, and TE2 that have

2. Results and discussion

To start with, we consider a three-layer optical slab waveguide with refractive indices of $n_1$, $n_2$, and $n_3$ (figure 1(a)). Media 1, and 3 form the cladding layers, and the medium 2 forms the core layer having a thickness of $t$. The waveguiding film supports a discrete number of waveguide modes that always lie below the air light-line. For transverse-electric (TE) polarized waveguide modes, they comply with the characteristic equation

$$\gamma_i t = \tan^{-1}(\gamma_1/\gamma_2) + \tan^{-1}(\gamma_3/\gamma_2) + m\pi,$$

with non-negative integer $m$ denoting the order of waveguide mode. The parameters $\gamma_1$, $\gamma_2$, and $\gamma_3$ are

$$\gamma_1 = k_0(\beta^2 - n_i^2)^{1/2}, \quad \gamma_2 = k_0(n_i^2 - \beta^2)^{1/2},$$

and

$$\gamma_3 = k_0(\beta^2 - n_i^2)^{1/2},$$

where $\beta$ is the wavevector of the guided modes and $k_0$ is the wavevector of photons in free space. Here, we use the phenethylammonium lead iodide (PEPI) perovskite, with chemical formula $(C_6H_5C_2H_4NH_3)_2PbI_4$, as the core of the waveguide. Its dielectric function can be artificially represented as [25]

$$\varepsilon_{\text{PEPI}}(E) = n_{\infty}^2 + \frac{A}{E^2 - E^2 + i\Gamma_X E}, \quad (1)$$

where $n_{\infty}$ is the refractive index of the passive medium, $A = 0.4 \text{ eV}^2$ is the oscillator strength of PEPI exciton, $E_X = 2.4 \text{ eV}$ is the energy of exciton resonance, and $\Gamma_X = 30 \text{ meV}$ is its linewidth. In the following, when taking the Lorentz term out, the structure is passive, and otherwise is active. The slab waveguide sits in an asymmetric dielectric environment with $n_1 = 1$ and $n_3 = 1.46$ respectively corresponding to air and silica.

We calculate in figure 1(b) the $E$–$\beta$ dispersion curves of TE-guided mode in the PEPI slab waveguide. The passive structure with $t = 300$ nm supports three guided modes TE0, TE1, and TE2 that have
Figure 1. (a) A three-layer dielectric slab waveguide with thickness of core layer $t$. The red, blue, and green shaded areas schematically represent 0th, 1st, and 2nd-order waveguide modes, respectively. (b) TE$_0$ ($m = 0$), TE$_1$ ($m = 1$), and TE$_2$ ($m = 2$) mode dispersion for the passive slab waveguide with $t = 300$ nm, where regions 1, 2, and 3 correspond to air ($n_1 = 1$), PEPI ($n_2 = \sqrt{\varepsilon_{\text{PEPI}}}$), and silica ($n_3 = 1.46$), respectively. The dotted curves indicate corresponding mode dispersion for an active structure. Also shown in the dispersion plot are two light lines: these are the air light-line, red-dashed, and the silica light-line, green-dashed. The golden solid curve is the imaginary component of the permittivity of PEPI perovskite. (c) Sketch of the perovskite-based PhC placed on top of a silica substrate, with the following fixed parameters: $a = 270$ nm, $h = 20$ nm, $w = 135$ nm. The inset illustrates the lattice structure of PEPI perovskite. (d) Folded quasi-TE guided mode dispersion of the passive PhC with a waveguide thickness $t = 300$ nm in the first Brillouin zone. The indices $(n, m)$ account for the diffraction order $(n)$ and the horizontal waveguide mode $(m)$. The dotted curves indicate the dispersion of $(1, 0)$ TE mode for an active structure, and the horizontal solid line, the position of the PEPI exciton resonance. The red-dashed line indicates the air light-line.

approximately linear dispersion with constant group velocity and asymptotically converge to the silica light-line at low energies. When the structure becomes active, each dispersion curve splitting two energy branches. And as the wavevector increases, the low energy branch approaches gradually the energy position at which the PEPI exciton has a maximum absorption band, as shown by the golden solid curve. Since these guided modes possess dispersive curves always lying below the air light-line, thus leading a momentum mismatch, they cannot be coupled to free-space photons, forbidding their far-field probing. By texturing the slab waveguide, as illustrated in figure 1(c), these pure bound states can be folded back to the first Brillouin zone via grating vectors. Due to the PhC having a shallow grating depth $h = 20$ nm, the scattering by the grating ridges and grooves is too weak to dramatically perturb the waveguide mode. In this situation, when the wavevector of TE-polarized incident light is rotated around the $y$ axis, the Bragg coupling condition can be expressed as $\beta = k_x + nG_x$, with $G_x = 2\pi/a$ being the reciprocal lattice vector and $n$ the diffraction order of the grating. The equation allows us to determine the $E-k_x$ dispersion relation of the $(n, m)$ TE-guided mode resonances. This is shown in figure 1(d), where $(\pm 1, 0), (\pm 1, 1), \text{ and } (\pm 1, 2)$ TE-guided modes are revealed for the passive PhC at a waveguide thickness $t = 300$ nm. These modes exhibit the folded dispersion beyond the air light-line (indicated as a red-dashed line), and therefore can be accessible in free space. When we turn the exciton resonance back on, a clear anticrossing of the $(1, 0)$ guided mode is
Figure 2. Transmittance spectra of the (a) and (e) passive and (b) and (f) active PhCs at different waveguide thickness $t$, which are calculated as a function of in-plane wavevector $k_x$ using FEM (grey colormap) and RCWA (thermal colormap) under TE-polarized incidence. (c) and (g) The fits to LP and UP dispersion from spectra in panels (b) and (f) with the coupled oscillator model, with Rabi splittings of $\Omega_{R} = 252.5$ and 258.6 meV, respectively. (d) and (h) Top: the associated Hopfield coefficients $|X|^2$ and $|C|^2$, denoting the exciton and photon fractions in the UP branch, respectively. Bottom: the lifetime of the UP, where the grey and green dashed lines denote the lifetime of uncoupled photon and exciton, respectively.

observed, as indicated by the dotted curves. The minimum energy splitting, that is, Rabi splitting, occurs where the frequency of the bare mode matches that of the PEPI exciton resonance (the horizontal solid line). Since the field profiles of the weakly corrugated PhC are very similar to those of the unpatterned waveguiding film, the Rabi energy for the quasi-bound modes is approximately amount to that for the pure bound ones.

Having established that resonance coupling between excitons and modes beyond the light-line may occur in the PEPI PhC, we begin to explore the behavior by accessing the far-field information of the device. For the sake of making our results convincing, we perform numerical calculations using two popular methods, FEM and RCWA (see methods for the calculation details). Figure 1(a) shows the calculated transmission map as a function of in-plane wavevector $k_x$ for a passive structure that has a thin waveguide thickness $t = 100$ nm in TE polarization, in which the two numerical algorithms exhibit greatly good agreement. The PhC structure supports two counterpropagating guided-mode resonances ($\pm 1, 0$) with linear dispersion. Owing the weak scattering from the shallower grating ridges and grooves, the resultant PhC modes exhibit Fano-like resonances with very high-quality factor. The formation of Fano-type resonance is due to the interference of the discrete waveguide mode with the continuum of vacuum states.

When the structure is switched to be active, as seen in figure 2(a), a broad, dispersionless resonance is shown up at $E = 2.4$ eV in the transmission map, which corresponds to the excitation of exciton mode in perovskite. As the guided modes approach the exciton energy, a clear anti-crossing dispersion is seen, which proves the formation of self-coupled upper (UP) and lower (LP) $X$–$C$ polariton modes, with $X$ and $C$ representing the exciton and the PhC cavity modes, respectively.

In order to determine the strong $X$–$C$ coupling, the anti-crossing can be fit to the coupled oscillator model given by

$$E_{UP,LP} = \frac{E_X + E_C}{2} + \frac{\Gamma_X + \Gamma_C}{2} \pm \sqrt{\left(\frac{1}{4}(E_X - E_C + \Gamma_X + \Gamma_C)^2\right)^2 + g^2 + \left(\frac{1}{4}(E_X - E_C + \Gamma_X + \Gamma_C)^2\right)^2},$$

where $E_X$ and $E_C$ are the energies of uncoupled PEPI exciton and PhC cavity modes, respectively; $\Gamma_X$ and $\Gamma_C$ are the half width half maximum of corresponding modes; $g$ is the coupling strength between exciton and
photon. At zero detuning, that is, $E_X = E_C$, the Rabi splitting can be obtained as $\Omega_R = \sqrt{4\epsilon^2 - (\Gamma_X - \Gamma_C)^2}$. In order to achieve the clear strong coupling regime, the splitting must satisfy the condition: $\Omega_R > \Gamma_X + \Gamma_C$. We fit the coupled oscillator model in equation (2) to the calculated data in figure 2(a), as seen figure 2(c). In the fitting, the exciton energy is fixed to 2.4 eV, while the uncoupled cavity mode is determined using a linear combination associated with $k_x$. The coupled dispersion branches fitted by equation (2) are shown by solid curves in figure 2(c), by which we obtain the coupling strength of $g = 126.7$ meV, and the yielding Rabi splitting of $\Omega_R = 252.5$ meV. Strong coupling occurs when $\Omega_R > \Gamma_X + \Gamma_C$, with $\Gamma_X = 28$ meV and $\Gamma_C = 6$ meV, which is satisfied for our system. Also, we give the fractions of exciton and photon in the UP branch (top panel, figure 2(d)) by calculating the Hopfield coefficient. The result implicates that the mixed polariton takes on a more photonic character away from the normal incidence, while it assumes a predominate excitonic behavior close to the normal incidence, to which the LP branch is exactly opposite. In addition, high quality-factor PhC modes allow us to engineer the exciton–polariton lifetime. This is shown by the bottom panel in figure 2(d), where the lifetime of the UP is calculated using the equation $\Gamma_{UP} = |X|^2\Gamma_X + |C|^2\Gamma_C$. One can note that the hybridized exciton–polariton has a dramatically enhanced lifetime compared to the bare exciton.

Same studies are carried out in the PEPI PhC at a waveguide thickness $t = 300$ nm, as seen by the bottom row in figure 2. In the system, one can achieve the strong coupling between higher-order PhC modes and the perovskite exciton (figure 2(i)), and obtains a Rabi energy of $\Omega_k = 258.6$ meV between exciton and degenerate ($\pm 1, 1$) guided modes. This enhanced coupling suggests that the guided mode at $t = 300$ nm has a better field confinement within the active material than the fundamental mode at $t = 100$ nm. And a thicker PEPI waveguiding film can give rise to the reduced lifetime of exciton, as shown by the bottom panel in figure 2(h). Nevertheless, it can be seen from this picture that high-order modes possess giant radiative lifetime, which thereby increases the lifetime threshold of the mixed light–matter state. In other words, the hybrid polaritons will inherit the lifetime of photonic modes at higher wavevectors.

In what follows, we give a comprehensive discussion of the effect of the waveguide thickness on the mode splitting. Figure 3(a) shows spectra-splitting energies of ($\pm 1, 0$) and ($\pm 1, 1$) modes at different waveguide thicknesses, as indicated by spheres. It is found that the Rabi splitting of the fundamental modes is larger than that of the higher-order ones, and that as the waveguide thickness increases the mode splitting is gradually grown. We can qualitatively shed light on this behavior by examining the electric field distribution within the slab waveguide. When the waveguide layer is thicker, the electric-field is mostly confined within the active material with low evanescent leakage into the external environment. Hence, this can effectively increase the volume of light–matter interaction, leading to the improved coupling. A quantitative insight into the capability of field confinement of each mode can be gained by evaluating the guided-mode width in a slab waveguide. Here, the effective mode width of a guided wave can be simply expressed as $W_{eff} = \left(\int_{-\infty}^{\infty} |E|^2 dz \right)^2 / \int_{-\infty}^{\infty} |E|^4 dz$, with $E$ being the electric field. And the field confinement is characterized using the expression $F = t/W_{eff}$. The calculated field confinement of ($\pm 1, 0$) and ($\pm 1, 1$) guided modes as a function of guiding-layer thickness $t$ is shown by the open triangles in figure 3(a). It can be seen that the variation has a same trend with that for the mode splitting, which illustrates that the vacuum electric field can serve as a main figure-of-merit for the PEPI PhC for strong coupling.

From the classical standpoint [28], the Rabi splitting is proportional to the square root of the volume density of emitters, that is, $\Omega_R \propto \sqrt{A} = \sqrt{N\epsilon^2/V\varepsilon_0 m_e}$, with $A$ being the oscillator strength, $N$ the number of excitons involved in the strong coupling per unit volume $V$, $\epsilon$ the elementary charge, $\varepsilon_0$ the permittivity in vacuum, and $m_e$ the effective electron mass. The oscillator strength is proportional to the exciton density, and thereby a stronger dipole moment will be generated with a larger density. Thus, the oscillator strength of the PEPI exciton is a second critical element controlling the strong coupling. Figure 3(b) displays the transmittance of the single-mode PEPI PhC ($t = 100$ nm) as a function of oscillator strength $A$. With increase of the oscillator strength, the exciton bandwidth becomes wider and the energy difference between LP and UP is substantially increased. Particularly, the yielding Rabi splitting reaches up to $\Omega_R = 375$ meV when $A = 0.85$. The value is much higher than that recently reported in perovskite-based metasurfaces, manifesting that guided-mode resonances constitute a flexible and efficient platform in the context of strong light–matter interaction.

Due to the in-plane anisotropy of one-dimensional dielectric gratings, the PEPI PhC can allow for more complex dispersion. When the incident electromagnetic wave with the magnetic vector parallel to the $x$ axis lies in the $yz$ plane, the expression for the momentum–energy dispersion of the guided modes can be written as

$$\beta = \sqrt{k_y^2 + \left(\frac{2\pi n}{a}\right)^2}$$

(3)
Figure 3. (a) Mode-splitting energy (spheres) and field confinement (open triangles) of the PhC modes at different waveguide thicknesses. (b) RCWA calculated transmission spectrum for the perovskite PhC with $t = 100$ nm as a function of coupled oscillator strength of exciton. All calculations are performed at zero detuning between guided-mode and exciton resonances.

Figure 4. The map of transmission of the (a) passive and (b) active PhCs with $t = 100$ nm as a function of in-plane wavevector $k_y$. In the left-hand half of the panel (b), the solid lines correspond to coupled oscillator model in equation (2), and the black and blue dashed lines to the bare exciton and guided modes, respectively. (c) Group velocity normalized to light speed in vacuum for the bare PhC mode and the mixed polaritons, LP and UP.

Unlike the folded waveguide dispersion we have discussed earlier, the equation gives rise to an approximately parabolic dispersion. As shown in figure 4(a), the dispersion plot for the passive PhC strongly resembles that of monolithic microcavities formed by distributed Bragg reflectors [10, 29], where the cavity mode manifests a low group velocity close to normal incidence. For the PEPI structure (figure 4(b)), the mode’s dispersion shows a clear anti-crossing band and the Rabi splitting of $\Omega_R = 253$ meV is obtained from the oscillator model, suggesting that the corresponding mode confinement is highly similar to that in the $k_x$ case.

For the sake of quantitatively evaluating these dispersions, we calculate the group velocity that is normalized to the light speed (that is, $|v_g|/c$) in free space from the analytical dispersion curves given by the coupled oscillator model. This is shown in figure 4(c), where the uncoupled PhC polariton, UP, and LP are highlighted by black, red, and blue spheres, respectively. Evidently, the index for the passive structure is linearly increased or decreased in a half wavevector space, while it shows bending for the PEPI PhC. More specifically, when operating out of the anti-crossing region (higher wavevector), the UP inherits the more cavity-mode characteristic exhibiting high group velocity, while the LP has a low group velocity due to its predominant exciton feature. And the LP experiences a slow-light region with a maximum group velocity $|v_g|/c \approx 0.063$ at anti-crossing points corresponding to $k_y \approx \pm 7.5$ rad $\mu$m$^{-1}$. In contrast to high-velocity regime ($|v_g|/c \approx 0.41$ as $k_x > 5$ rad $\mu$m$^{-1}$, figure 2(b)), it may be highly favorable to observation of some intriguing effects in the slow-light regime where high density of states is formed, such as nonlinear effects and Bose–Einstein condensation of polaritons.

We finally provide the discussion with respect to the strong coupling of the PhC modes with PEPI excitons under transverse-magnetic (TM) polarized light. As seen in figure 5(a), the appearance of strong coupling is also clearly observed, with Rabi-splitting of $\Omega_R = 203$ meV, which, nevertheless, cannot be achieved in PhCs coupled to two-dimensional exciton platforms. For example, excitons in TMDC
monolayers interacting with TM-polarized modes in PhCs have been confirmed to be in the weak-coupling regime [30]. This is because excitons in monolayers have the polarization of the transition dipole moment predominantly in the in-plane direction. More importantly, the photonic cavities produce a weak in-plane electric field in the TM polarization, thus leading to a very weak coupling with in-plane excitons in monolayers. In order to elucidate this point, we carry out numerical simulations of the near-field response of the passive PhC employing FEM. This is displayed in figure 5(b), where the electric-field profiles in the in-plane ($E_x$) and out-of-plane ($E_z$) directions are plotted for the bare TM-polarized PhC mode at zero detuning. It can be seen that the enhancement factors for $E_x$ and $E_z$ fields respectively are $\sim 3$ and $\sim 5$. Such weak in-plane electric field is insufficient to produce strong interaction with the electronic transition of a monolayer. But this can be resolved for bulk exciton media due to the presence of out-of-plane excitons together with strong out-of-plane electric field.

3. Conclusion

In conclusion, we have theoretically demonstrated a subwavelength photonic system made of PEPI perovskite, in which strong coupling between excitons with guided-mode resonances occurs within the same nanophotonic object. Our results showed that the giant Rabi splitting is obtained for the system excited by either TE or TM polarizations owing to the large localization of electromagnetic field inside exciton materials. This implies that the photonic system supporting guided-mode resonances could be a useful playground for studying the self-hybridization of excitons and for exploring a rich diversity of polaritonic phenomena. In addition, we demonstrated that the dispersions and the lifetime of exciton–polaritons can be engineered via propagating guided-mode polaritons, thus enabling the realization of moving exciton–polariton condensates. More generally, one would like to emphasize that the self-hybridized exciton–polaritons have been recently observed in several different platforms based on nanostructured perovskite dielectrics, including nanowires [31, 32], isolated nanoparticles [33], hybrid plasmonic–photonic systems [34] and metasurfaces [23]. Complementary to these studies, here we opened up a new adventure to exploring strong light–matter interactions using perovskite PhCs.

Methods. The RCWA simulations of angular-resolved transmittance have been performed with S$^4$, a freely available software package provided by Fan Group in the Stanford Electrical Engineering Department [35]. Electromagnetic and transmittance simulations have been done using a commercial software based on FEM (COMSOL Multiphysics). Our simulations are three-dimensional with the Bloch boundary conditions in the $x$ and $y$ directions and perfectly matched layers along the $z$ direction to ensure absorption of outgoing radiated waves. The structure is illuminated with a broadband plane-wave light source incident and an incident angle $\theta$, and the projection of incident wavevector onto the structural surface along the $x$ and $y$ directions respectively are $k_x = \frac{2\pi}{\lambda} \sin(\theta) x$ and $k_y = \frac{2\pi}{\lambda} \sin(\theta) y$. Electric field direction is determined based on the required polarization to provide TM across the grating or TE along the grating wires.
Acknowledgments

This work is supported by the National Natural Science Foundation of China (62105095, U1804261), the Natural Science Foundation of Henan Province (202300410238), and the National Scientific Research Project Cultivation Fund of Henan Normal University (20210381, 20211P22).

Data availability statement

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

ORCID iDs

Lixia Li https://orcid.org/0000-0003-4830-7899
Wanlu Li https://orcid.org/0000-0003-2538-5894
Yufang Liu https://orcid.org/0000-0003-1896-8864

References

[1] Kasprzak J et al 2006 Bose–Einstein condensation of exciton polaritons Nature 443 409–14
[2] Amo A, Lefrère J, Pigeon S, Adrados C, Ciuti C, Carusotto I, Houdré R, Giacobino E and Bramati A 2009 Superfluidity of polaritons in semiconductor microcavities New J. Phys. 11 043029
[3] Christopoulos S et al 2007 Room-temperature polariton lasing in semiconductor microcavities Phys. Rev. Lett. 98 126405
[4] Nikolis V C et al 2019 Strong light–matter coupling for reduced photon energy losses in organic photovoltaics Nat. Commun. 10 3706
[5] Wang Y, Knoll W and Dostalek J 2012 Bacterial pathogen surface plasmon resonance biosensor advanced by long range surface plasmons and magnetic nanoparticle assays Anal. Chem. 84 8343–50
[6] Chen J, Nie H, Peng C, Qi S, Tang C, Zhang Y, Wang L and Park G-S 2018 Enhancing the magnetic plasmon resonance of three-dimensional metal optical metasurfaces via strong coupling for high-sensitivity sensing J. Lightwave Technol. 36 3481–5
[7] Chen J, Yang C, Gu P, Kuang Y, Tang C, Chen S and Liu Z 2021 High sensing properties of magnetic plasmon resonance by strong coupling in three-dimensional metasurfaces J. Lightwave Technol. 39 562–5
[8] Ciers J, Roch J G, Carlin J-F, Jacopin G, Butté R and Grandjean N 2017 Propagating polaritons in III-nitride slab waveguides Phys. Rev. Appl. 7 034019
[9] Zheng D, Zhang S, Deng Q, Kang M, Nordlander P and Xu H 2017 Manipulating coherent plasmon–exciton interaction in a single silver nanorod on monolayer WSe2 Nano Lett. 17 3809–14
[10] Liu X, Galšky T, Sun Z, Xia F, Lin E-C, Lee Y-H, Kema–Cohen S and Menon V M 2019 Strong light–matter coupling in two-dimensional atomic crystals Nat. Photon. 9 30–4
[11] Wang S, Mika A, Hutchison J A, Genet C, Jouaiit A, Hosseini M W and Ebbesen T W 2018 Phase transition of a perovskite strongly coupled to the vacuum field Nanoscale 6 7243–8
[12] Zakharko Y, Graf A and Zaumseil J 2016 Plasmonic crystals for strong light–matter coupling in carbon nanotubes Nano Lett. 16 6504–10
[13] Huang Y, Wu F and Yu L 2020 Rabi oscillation study of strong coupling in a plasmonic nanocavity New J. Phys. 22 063053
[14] Hou L, Wang Q, Zhang H, Wang P, Gan X, Xiao F and Zhao J 2022 Simultaneous control of plasmon–exciton and plasmon–trion couplings in an Au nanosphere and monolayer WS2 hybrid system APL Photon. 7 02617
[15] Conteduca D, Reardon C, Scullion M G, Dell’olio F, Armenise M N, Krauss T F and Ciminelli C 2017 Ultra-high Q/V hybrid cavity for strong light–matter interaction J. Lightwave Technol. 39 562–5
[16] Wang Q, Sun L, Zhang B, Chen C, Shen X and Lu W 2016 Direct observation of strong light-exciton coupling in thin WS2 flakes Opt. Express 24 7151–7
[17] Munkhbat B, Baranov D G, Stührenberg M, Wersäll M, Binnerts D and Shegai T 2019 Self-hybridized exciton–polaritons in multilayers of transition metal dichalcogenides for efficient light absorption ACS Photon. 6 139–47
[18] Zhang X, De-Eknamkul S, Sanvitto D, Liew T C H and Xiong Q 2021 Perovskite Weyl semimetal WTe2 Nano Lett. 21 1315–24
[19] Zheng D, Zhang S, Deng Q, Kang M, Nordlander P and Xu H 2017 Manipulating coherent plasmon–exciton interaction in a single silver nanorod on monolayer WS2 Nano Lett. 17 3809–14
[20] Zhang H, Abhiraman B, Zhang Q, Miao J, Jo K, Roccaesca S, Knight M W, Davoyan A R and Jariwala D 2020 Hybrid exciton–plasmon–polaritons in van der Waals semiconductor gratings Nat. Commun. 11 3552
[21] Zhang X, Zhang H, Kuang Y, Wang L, Zhao M, Charlie Johnson A T, Tongay S and Cubukcu E 2020 Ultra-thin WS2 on-glass photonic crystal for self-resonant exciton–polaritons Adv. Opt. Mater 8 1901988
[22] Fieramosca A et al 2019 Two-dimensional hybrid perovskites sustaining strong polariton interactions at room temperature Sci. Adv. 5 eaax9967
[23] Su R, Fieramosca A, Zhang Q, Nguyen H S, Deleporte E, Chen Z, Samwito D, Liew T C H and Xiong Q 2021 Perovskite semiconductors for room-temperature exciton–polaritons Nat. Mater. 20 1315–24
[24] Su R, Diederichs C, Wang J, Liew T C H, Zhao J, Liu S, Xu W, Chen Z and Xiong Q 2017 Room-temperature polariton lasing in all-inorganic perovskite nanotetraplates Nano Lett. 17 3982–8
[25] Kang M, Gerace D, Drouard E, Trippe-Allard G, Léée F, Mazurczyk R, Deleporte E, Seassal C and Nguyen H S 2020 Tailoring dispersion of room-temperature exciton–polaritons with perovskite-based subwavelength metasurfaces Nano Lett. 20 2113–9
[26] Tan C, Yue Z, Dai Z, Bao Q, Wang X, Lu H and Wang L 2018 Nanograting-assisted generation of surface plasmon polaritons in Weyl semimetal WTe2 Opt. Mater. 66 421–3
[27] Lu H et al 2020 Magnetic plasmon resonances in nanostructured topological insulators for strongly enhanced light–MoS₂ interactions Light Sci. Appl. 9 191
[28] Törnä P and Barnes W L 2015 Strong coupling between surface plasmon polaritons and emitters: a review Rep. Prog. Phys. 78 013901
[29] Liu X, Bao W, Li Q, Ropp C, Wang Y and Zhang X 2017 Control of coherently coupled exciton polaritons in monolayer tungsten disulphide Phys. Rev. Lett. 119 027403
[30] Zhang L, Gogna R, Burg W, Tutuc E and Deng H 2018 Photonic-crystal exciton–polaritons in monolayer semiconductors Nat. Commun. 9 713
[31] Shang Q et al 2018 Surface plasmon enhanced strong exciton–photon coupling in hybrid inorganic–organic perovskite nanowires Nano Lett. 18 3335–43
[32] Zhang S et al 2018 Strong exciton–photon coupling in hybrid inorganic–organic perovskite micro/nanowires Adv. Opt. Mater. 6 1701032
[33] Tiguntseva E Y et al 2018 Tunable hybrid Fano resonances in halide perovskite nanoparticles Nano Lett. 18 5522–9
[34] Lu L, Le-Van Q, Ferrier L, Drouard E, Seassal C and Nguyen H S 2020 Engineering a light–matter strong coupling regime in perovskite-based plasmonic metasurface: quasi-bound state in the continuum and exceptional points Photom. Res. 8 A91–100
[35] Liu V and Fan S 2012 S4: a free electromagnetic solver for layered periodic structures Comput. Phys. Commun. 183 2233–44