Spatially coherent surface resonance states derived from magnetic resonances

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Abstract. A thin metamaterial slab comprising a dielectric spacer sandwiched between a metallic grating and a ground plane is shown to possess spatially coherent surface resonance states that span a large frequency range and can be tuned by structural and material parameters. They give rise to nearly perfect angle-selective absorption and thus exhibit directional thermal emissivity. Direct numerical simulations show that the metamaterial slab supports spatially coherent thermal emission in a wide frequency range that is robust against structural disorder.

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1. Introduction

Surface plasmon polaritons (SPPs) can modulate light waves at the metal–dielectric interface with wavelengths much smaller than that in free space [1], which enables the control of light in a subwavelength scale for nanophotonic devices [2]. SPPs with large coherent lengths are useful in many areas, including optical processing, quantum information [3] and novel light–matter interactions [4]. The enhancement of local fields by SPPs is particularly important as it opens a new route to absorption enhancement [5], nonlinear optical amplification [6, 7] and weak signal probing [8, 9]. As the properties of SPP are pretty much determined by the natural (plasmon) resonance frequency, there is not much room for us to adjust the SPP response for practical applications. With induced surface current oscillations on an array of metallic building blocks [10]–[16], metamaterial surfaces can manipulate electromagnetic waves in a similar way to SPPs. Such SPPs or surface resonance states on structured metallic surfaces are tunable by geometric parameters.

In this paper, we examine the properties of surface resonance states at a dielectric–metamaterial interface that exhibits a magnetic response to the incident waves and strong local field enhancement. We will see that these surface resonance states can give highly directional absorptivity and emissivity, and may thus help us to realize interesting effects, such as spatially coherent thermal emission. As the structure is very simple, it can be fabricated down to the infrared (IR) and optical regime [17]–[19].

We will show that a thin metamaterial slab, with a thickness much smaller than the operational wavelength, supports delocalized magnetic surface resonance states with a long coherent length in a wide range of frequencies. Operating in a broad frequency range, these spatially coherent SPPs are surface resonance states with quasi-transverse magnetic (TEM) modes guided in the dielectric layer that are weakly coupled to free space, and the coupling strength can be controlled by tuning structural parameters, while the frequency can be controlled by varying the structural and material parameters. The high fidelity of these surface resonance states results in directional absorptivity or emissivity, which is angle dependent with respect to frequency. Finite-difference-in-time-domain (FDTD) simulations verify that the highly directional emissivity from the slab persists in the presence of structural disorder in the grating layer.

Such metal–dielectric–metal (MDM) structures were recognized as artificial magnetic surfaces with high impedance by the end of the last century [12]. The incident waves induce surface current solenoids on the unit cells of the ultrathin MDM structure, giving rise
to magnetic induction, which can be described with an effective permeability in Lorentz type [12, 20]. After the concept of metamaterial was proposed [21], A P Hibbins and co-workers numerically and experimentally proved that ultrathin MDM structures can resonantly absorb or transmit radiations in the low frequency limit [22]. They proposed that the central frequencies of absorption peaks are independent of the incident angle with an interpretation of the Fabry–Perot resonant mode (equation (1)) in [22]). The same group further explored the angle-independent absorption as the main scenario of the incremental work by measuring the flat bands of surface wave dispersion in the visible [23] and the microwave regions [24]. In contrast, we found that the structures with proper design also support very narrow absorption peaks that are sensitive to the incident angle and obviously do not satisfy the Fabry–Perot resonance condition suggested in the previous studies.

It is worth noting that an angle-independent peak is quite different from an angle-dependent one in physical origin. The former, investigated in [22]–[24], comes from the localized surface resonance states, while the latter, found by us, comes from the collective surface resonance states. Mode analysis presents an intuitive picture of the formation of these collective surface states. When high-order quasi-TEM modes are the dominant components of the guided waves inside the dielectric layer, phase correlation is assigned to the outgoing waves emitted from the air slits of the grating, giving rise to a collective response. And the spatial coherence of surface resonant states will survive, provided that both the leakage from the dielectric layer to air slits and the material absorption are weak enough. As the interaction between the structure and the incident waves will excite quasi-TEM modes inside the dielectric layer, the magnetic induction must be parallel to the MDM surfaces if it exists. Thus, a surface resonance state on an MDM structure is usually magnetic in nature. Our findings about spatially coherent surface resonance states are original compared to the common knowledge, and have great potential in the coherent control of SPPs and thermal emission radiations.

2. Model description and mode expansion method

Our model system is schematically illustrated in figure 1. Lying on the \( \hat{x}\hat{y} \)-plane, the slab comprises an upper layer of a metallic lamellar grating with thickness \( t \), a dielectric spacer layer as a slab waveguide with thickness \( h \) and a metallic ground plane. The metallic strips are separated by a small air gap \( g \), giving rise to a period of \( p = a + g \) for the lamellar grating.

Figure 1. Schematic diagram of the magnetic metamaterial slab. The geometric parameters are \( t = 0.2 \mu m \), \( h = 0.8 \mu m \), \( a = 3.8 \mu m \), \( g = 0.2 \mu m \) and \( p = a + g = 4.0 \mu m \). The dielectric spacer layer is slightly dissipative by assigning a complex permittivity \( \varepsilon_{III} = \varepsilon_r \varepsilon_0 + i (\sigma/\omega) \) with \( \varepsilon_r = 2.2 \) and \( \sigma = 66.93 \text{ S m}^{-1} \).
The geometric parameters of our model are \( t = 0.2 \mu m, h = 0.8 \mu m, a = 3.8 \mu m, g = 0.2 \mu m \) and \( p = a + g = 4.0 \mu m \). Each metallic strip, together with the ground plane beneath it, constitutes a planar resonant cavity as the building block that gives magnetic responses at cavity resonances [12, 16]. As the metallic grating is along the \( \hat{x} \)-direction, the guided waves in the dielectric layer (at \( 0 < z < h \) in region III) will always couple to the incident waves with a non-zero component \( E_x \neq 0 \) of the electric field.

As a first step, we consider a transverse magnetic (TM) polarized incident plane wave in the free semi-space (at \( z > h + t \) in region I). The electric field \( \vec{E} \) lies in the \( \hat{x}\hat{z} \)-plane, the magnetic field \( \vec{H} \) is along the \( \hat{y} \)-axis and the in-plane wave vector is \( \vec{k}_0 = k_x \hat{e}_x \) (\( k_y = 0 \)). The total magnetic fields in region I and in region III can be written in terms of the reflection coefficients \( r_m \) and the guided Bloch wave coefficients \( t_m \) [25]–[28], as

\[
\begin{align*}
H_y^I(\vec{r},z) &= \delta_{m,0} e^{-ik_m^I z} e^{ik_m^I x} + \sum_{m} r_m e^{ik_m^I z} e^{i(k_m^I + 2\pi n/p)x}, \\
H_y^II(\vec{r},z) &= \sum_{m} t_m e^{ik_m^II z} e^{-ik_m^II (z-h)} e^{i(k_m^I + 2\pi n/p)x},
\end{align*}
\]

where the term \( \delta_{m,0} e^{-ik_m^I z} e^{ik_m^I x} \) denotes the incident plane wave with \( \delta_{m,0} \) being the Kronecker function and \( m \) being the Bloch order; \( e^{i(k_m^I + 2\pi n/p)x} \) denotes the wave component of the \( m \)th Bloch eigenmode in the semi-free space (region I) and the dielectric layer (region III) with respect to \( \vec{k}_m = \hat{x}(k_x + 2m\pi/p) \). \( \vec{k}_m \) is the in-plane wave vector and \( \hat{G}_m = \hat{x} \cdot 2\pi m/p \) is the \( m \)th reciprocal lattice vector. \( k_m^I = \sqrt{\varepsilon_0 \mu_0 \omega^2 - |\vec{k}_m|^2} \) and \( k_m^II = \sqrt{\varepsilon_{III} \mu_0 \omega^2 - |\vec{g}_m|^2} \) are the \( z \)-components of the wave vector for the \( m \)th order Bloch eigenmode in region I and region III, respectively. \( \varepsilon_0 \) and \( \varepsilon_{III} \) are the permittivities of the vacuum and the dielectric, and \( \mu_0 \) is the vacuum permeability.

We shall mainly consider IR frequencies, at which the metals can be well approximated as perfectly electric conductors (PEC). The EM fields at \( h \leq z \leq h + t \) in region II are squeezed inside the air gaps, in which the magnetic fields can be expressed in terms of the expansion coefficients \( a_l \) and \( b_l \) of forward and backward guided waves, as

\[
H_y^II(\vec{r},z) = \sum_l [a_l e^{-i\beta_l(z-h+t)} + b_l e^{i\beta_l(z-h)}] g_l(x),
\]

where \( g_l(x) = \cos(l\pi/g(x+g/2)) \) (\( l = 0, 1, \ldots, n, \ldots \)) is the in-plane distribution of the \( l \)th guided mode \( \alpha_l \) running over all air gaps [29]. \( \beta_l = \sqrt{\varepsilon_0 \mu_0 \omega^2 - (l\pi/g)^2} \) is the \( z \)-component of the wave vector for the \( l \)th guided mode \( \alpha_l \).

We can obtain the coefficients \( t_m(f, \vec{k}_0) \) and \( r_m(f, \vec{k}_0) \) of the \( m \)th guided and reflected waves by applying the boundary continuity conditions for the tangential components of electromagnetic wave fields (over the slits) at the interfaces \( z = h \) and \( z = h + t \). Given that surface resonance modes are an intrinsic response, we can also assign zero to the incident plane wave and apply the boundary continuity conditions to the tangential components of wave fields to derive the eigenvalue equations. A surface resonance state can be determined by searching a zero value/minimum of the eigenvalue determinant in the reciprocal space provided that it is non-radiative/radiative with infinite/finite lifetime below/above the light line in free space. In general, we also developed the method for any specific wavevector and any specific polarization of an incident plane wave.

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Figure 2. Absorption spectra under TM-polarized incidence (a) as a function of frequency at incident angles of $\theta = 0^\circ$, $5^\circ$, $20^\circ$ and $30^\circ$ ($g = 0.2$ $\mu$m) and (b) as a function of incident angle at 54.3 THz (solid line, $g = 0.2$ $\mu$m), 40 THz (dashed line, $g = 0.2$ $\mu$m) and 40 THz (dotted line, $g = 0.1$ $\mu$m).

3. Absorption spectra properties and spatial coherence of magnetic surface resonance states

We derived the absorption spectra of the slab $A(\vec{k}_0, \omega) = 1 - \sum_m \text{Re}(k_{m\omega}^I/k_{\omega 0}) r_m(\vec{k}_0, \omega)|^2$, which includes the contributions from all Bloch orders of reflected waves. As a consequence, $A(\vec{k}_0, \omega)$ gives information about the surface resonance states as well as the emissivity properties as governed by Kirchhoff’s law [30]. We shall assume that the dielectric spacer layer is slightly dissipative by assigning a complex permittivity $\varepsilon_{\text{III}} = \varepsilon_r \varepsilon_0 + i(\sigma/\omega)$ with $\varepsilon_r = 2.2$ and $\sigma = 66.93$ $\text{S m}^{-1}$ [$\text{Im}(\varepsilon_{\text{III}}) \approx 10^{-2}\varepsilon_r \varepsilon_0$] in the calculated frequency regime. In figure 2(a), we present the absorption spectra at various incident angles. The spectra exhibit a low, broad peak at 13.2 THz, which is almost independent of the incident angle, while the other absorption peaks at higher frequencies are narrow and sensitive to the incident angle with a maximum absorption approaching 100%. The slab thus acts as an all-angle absorber at 13.2 THz (a similar result can be found in [31]) but exhibits sharp angle-selective absorption peaks at higher frequencies. Shown as solid and dashed lines in figure 2(b), the sharp angular dependence of absorption coefficients (note that the vertical axis is in log-scale) at 40.0 and 54.3 THz implicitly implies the existence of spatially coherent surface resonance states. The angle-dependent absorption peaks become lower and disappear gradually with an increase in the material loss. This presents a way to realize nearly perfect absorption with weakly absorptive materials by coherent surface resonance states. The coherent length of a surface resonance state can be estimated by the ratio of the wavelength $\lambda$ and the full-width at half-maximum (FWHM) $\Delta \theta$ of the absorption.
peak [32]. For example, for the $\Gamma_4$ state at 54.3 THz and $\vec{k}_0 = 0$, the angular FWHM of the corresponding absorption peak $\Delta \theta = 4.6^\circ$ (from $\theta = -2.3^\circ$ to $\theta = 2.3^\circ$) gives rise to a coherent length $\lambda / \Delta \theta = 68.5 \mu m \approx 12.4 \lambda$. The coherent length is about 220$\lambda$ for the surface resonance state at 50.22 THz and $\vec{k}_0 = 0.01 \pi \cdot p$ with $\Delta \theta = 0.26^\circ$ (not shown in figure). The angular FWHM is reduced if the gap size is smaller, as shown by the dashed and dotted lines in figure 2(b) for $g = 0.2$ and 0.1 $\mu m$ at 40 THz, which means that the coherent length of the surface resonant modes can be controlled by the gap period ratio $g / p$. It is worth noting that, although $k_y = 0$ is assumed for the calculated results shown in figure 2, the angle-dependent absorption peaks are readily obtained for any specific incident angle.

To quantitatively characterize the formation of these spatially coherent surface resonance states, we employ the eigenmode expansion method to calculate the surface resonance dispersion (in the limit of no material loss), as shown in figure 3(b). The $B_1$ surface resonance states lie below the light line $L_2$ (magenta dashed line) and thus are non-radiative as evanescent modes. The surface resonances labeled as $B_2$ originate from the coupling of the fundamental magnetic resonance modes of the metal strip structure with the free space light line $L_2$. The surface resonances $B_3$ and $B_4$ are harmonic modes of the magnetic resonances that hybridize with the guided mode inside the dielectric layer. The calculated reflection phase difference between the 0th order reflected and incident electric fields, as shown in figure 3(a) for normal incidence (red line) and 2$^\circ$ incidence (blue line), clearly shows that the resonances are magnetic in nature when the surface resonances intersect the zone center at $\Gamma_2$ (13.2 THz) and $\Gamma_4$ (54.3 THz), as the reflection phase is zero, just like what a magnetic conductor surface does to the incident waves. The state $\Gamma_3$, invisible in the reflection phase spectrum under normal incidence (red solid line in figure 3(a)), is a dark state as its eigenmode is in mirror symmetry about the $yz$-plane and cannot couple with free space photons, while the other $B_3$ states

Figure 3. (a) Dispersion diagram of TM-polarized surface resonance states. (b) The reflection phase difference between the 0th order reflection and the TM-polarized incident plane wave at incident angles of $\theta = 0^\circ$ (red line) and $\theta = 2^\circ$ (blue line).
Figure 4. Spatial distributions of magnetic fields and electric fields in the $\hat{x}\hat{z}$-plane for $\Gamma_2$ state at $f_{\Gamma_2} = 13.2$ THz, $\theta = 0^\circ$ (a, c), a state on $B_3$ at $f = 50.22$ THz, $\theta = 2^\circ$ (b, e) and $\Gamma_4$ state at $f_{\Gamma_4} = 54.3$ THz, $\theta = 0^\circ$ (c, f).

can couple with external light under oblique incidence (see the blue line in figure 3(a)). For example, there exists in-phase reflection at a frequency of 50.22 THz under an incident angle of $2^\circ$, corresponding to a $B_3$ state at a frequency of 50.22 THz and $k_0 \parallel \approx 0.02 \pi / p$.

The angle-independent absorption peak at 13.2 THz is due to the $B_2$ mode, which is only weakly dispersive near the zone center. The more dispersive $B_3$ and $B_4$ modes are accountable for the incident-angle-sensitive absorption at the higher frequencies in figure 2(a). The field patterns in figures 4(a)–(c) present the spatial distribution of the real part of magnetic fields excited by the incident plane waves with incident angles $0^\circ$, $2^\circ$ and $0^\circ$ for the three surface resonance states on $B_2$, $B_3$ and $B_4$, respectively, and the corresponding vector diagrams of electric fields are shown in figures 4(d)–(f). We can see clearly that the electric fields reach a maximum in strength at the slab upper surface, and exponentially decay along the surface normal into the free space. This is precisely a picture of SPP modes. The field patterns come from the coincidence of the evanescent wave components at high Bloch orders at both sides of the metallic grating.

Although only TEM-guided modes are allowed to be excited in the thin MDM slab within the frequency of our interest, the $B_2$ states are quite different from the $B_3$ and $B_4$ states in field patterns inside the dielectric layer. We see from figures 4(b), (c), (e) and (f) that for a $B_3$ or $B_4$ state, there are nodes and anti-nodes in field patterns, while for the $B_2$ state, the magnetic field is almost uniformly distributed. Calculations on local field enhancement inside the dielectric slab resolve the puzzle. The black solid line in figure 5 presents the normalized magnetic field $|H|$ inside the dielectric with respect to that of incidence $|H_0|$ under an incident angle of $\theta = 2^\circ$. The Fourier component in $m = 0$ order (blue solid line) contributes the most at 13.2 THz and the least at 50.22 and 54.3 THz, while it is just the opposite for the contributions in combination from the two high-order Fourier components with $m = \pm 1$ (red solid line). Figure 5 also indicates that the enhancement of local field of an excited $B_3$ or $B_4$ state can be ten times larger than that of an excited $B_2$ state, the enhancement factor at 50.22 THz is about 100 times, while it is only 10 times at 13.2 THz. Thus, figures 4 and 5 clearly depict the physical origins of the two different surface resonance states.
Figure 5. Magnetic field \( \vec{H} \) inside the dielectric layer normalized to that of incidence \( H_0 \) under an incident angle of \( \theta = 2^\circ \). Black line: all the Bloch orders of TEM-guided modes included; red line: only the 0th Bloch order considered; blue line: summation of \(-1\)st and \(+1\)st Bloch orders of TEM-guided modes.

We see from figure 3(b) that the surface resonance dispersion of the slab comes from the interaction between the magnetic resonances and the (folded) light lines \( L_1 \) (for dielectrics) and \( L_2 \) (for air) grazing on the interfaces. In the limit of a small gap period ratio \((g/p = 0.05\) for example), our system is weakly Bragg-scattered, and as such, when a surface resonance state on branch \( B_3 \) or \( B_4 \) is excited, the induced wave fields inside the dielectric of region III are guided quasi-TEM modes dominated by \( \pm 1\)st Bloch orders. For that reason, the \( B_3 \) and \( B_4 \) states have high fidelity even though they are leaky modes, like most of their Bloch wavefunction components lying outside the free space light line. As the air gaps of the metallic grating serve to couple the electromagnetic waves of region I and region III, the quality factor of a resonance state can be estimated with the overlap integral of the tangential electric fields between the fundamental waveguide mode \( \{g_{l=0}(x) = \cos[l\pi g(x + g/2)] = 1\} \) in the air gap and the dominant Bloch waves \( \vec{E}_m^i \) \((i = I, III)\) in region I or region III for the coupling coefficients

\[
C_m^i = \frac{1}{w_x} \int_0^p \frac{k_{z m}^i}{\sqrt{\varepsilon_r k_x^2 + k_{z m}^2}} e^{-i(k_x + G_m)x} g_0(x)dx
\]

\[
= \frac{1}{\sqrt{\varepsilon_r k_x^2 + k_{z m}^2}} \sin c \left[ \frac{(k_x + G_m)g}{2} \right],
\]

when the air gap width \( g \ll p \) is satisfied. For the \( B_2 \) states, the major Fourier component of the wavefunction is \( \vec{E}_m^{III} \) in zero order, and as the corresponding \( z \) wavevector component \( k_{z0}^{III} \) in the dielectric layer (region III) is generally not small, \( C_0^{III} \) is usually very large according to equation (3), and the \( B_2 \) states leak out easily. The states on branches \( B_3 \) and \( B_4 \) have major Fourier components in \( m = \pm 1 \) order, and as they are asymptotic to the (folded) dielectric light lines \( L_1 \), the absolute value of \( k_{z m}^{III} \) \((m = +1\) for \( k_x < 0 \) or \( m = -1\) for \( k_x > 0 \)) is very small, resulting in the small coupling coefficients \( C_{-1}^{III} \) or \( C_{+1}^{III} \). The \( B_3 \) and \( B_4 \) modes have to travel a long distance before they leak out. They have a long lifetime and good spatial coherence. This also explains why the state \( \Gamma_3 \), a state precisely superposing on folded light line \( L_1' \) in the dielectric layer, is dark to the incident plane wave as \( k_{z m}^{III} = 0 \).
Different from $B_3$ and $B_4$ states, the $B_2$ states have a major Fourier component in $m = 0$ order, which directly couples to the free space photons. As a consequence, the $B_2$ states, forming a flat band far away from the light line $L_2$ when $k_0$ is small, are localized with resonant frequency scaled by the local geometry of the unit cell. The high mode fidelity of a $B_3$ or $B_4$ state also gives rise to a much more intense local field compared to the $B_2$ states. As shown in figure 5, the induced local field is 100 times stronger than the incident field for the state on $B_3$, while it is only 10 times stronger for $\Gamma_2$, and this is consistent with the absorptivity shown in figure 2(a). In addition, the coherent length can be adjusted by the gap width as the kernel $C_m^l$ is proportional to the gap period ratio $g/p$. Further calculations demonstrate that the angular FWHM of the absorption peak is reduced from 0.26° to 0.16° when the gap is decreased from 0.2 to 0.1 $\mu$m, corresponding to a coherent length of 358$\lambda$.

We note that most of the attention in previous studies has been devoted to the localized $B_2$ states [22]–[24], [31]. The spatially coherent surface resonance states will provide us with a new solution to coherent control of thermal emission radiations. J-J Greffet and co-workers showed that highly directional and spatially coherent thermal emission can be obtained by etching a periodic grating structure into an SiC surface [32]–[35]. The magnetic resonant modes in our system can do the same, as will be demonstrated below. Our system has the advantage that the operational frequency is tunable by changing the structural parameters, and the operational bandwidth is wide. In addition, our structure supports all-angle functionality for some specific range of frequencies, as shown in figure 6, although it is periodic only in one direction.

4. Coherent thermal emission

We performed FDTD simulations to emulate the emissions from a slab containing point sources with random phases using the same configuration parameters as those mentioned above. We purposely put disorder in the structure to test the robustness of the phenomena. We assigned two Gaussian distributions (they can be uniform distributions or other types as well) independently to the width of metallic strips and the center positions of air gaps to introduce a 4% (standard deviation) structural disorder. The slab has a lateral size of 60 periods along the $\hat{x}$-direction. A total of 1200 point sources with random phases are placed at the mesh points inside the dielectric layer. Directional emissions of a wide range of frequencies above 34 THz are confirmed by the simulation. The 4% structural disorder has little impact on the directional emissivity. Figures 6(a), (c) and (e) show the far-field emission patterns in the $\hat{x}\hat{z}$-plane ($H$-plane) at 40.0, 54.3 and 58.0 THz. The inset of figure 6(c) is a control calculation in which the top metal gratings are removed, so that there is just a dielectric layer with random phase sources above a metal ground plane. The directivity of emission from the random sources is lost. Figures 6(b), (d) and (e) present the absorptivity (under plane wave incidence) as a function of the in-plane wavevector (solid angle) at these frequencies. The results in figures 6(b), (d) and (f), calculated by the modal expansion theory for the $E_y = 0$ incidence, are in good agreement with the FDTD simulations shown in figures 6(a), (c) and (e). The strong angle selectivity of the absorption is evident and, by Kirchhoff’s law, the thermal emission should also be highly directional, which is a direct consequence of the good spatial coherence of the surface resonance states. As shown in figures 6(b), (d) and (f), the absorption/emission peaks generally trace out an arc in the $k_x$–$k_y$ plane but, near 54.3 THz (figure 6(d)), the dominant emission beam is restricted to a small region near the zone center. This is because the $\Gamma_4$ state is a minimum point if we
Figure 6. Radiation patterns in the $H$-plane (calculated by FDTD) and absorptivity (calculated by the mode expansion method) as a function of in-plane wavevectors at $f = 40.0 \text{THz}$ (a, b), $f = 54.3 \text{THz}$ (c, d) and $f = 58.0 \text{THz}$ (e, f). In FDTD simulations, 1200 point sources with random phases are placed at the mesh points inside the dielectric layer. A 4% structural disorder is included in the 80 $\mu$m simulation cell, which accounts for the slight asymmetry of the radiation patterns, but also demonstrates the robustness of the angle selectivity with respect to disorder. The inset of (c) shows a control calculation in which the top metal gratings are removed, so that there is just a dielectric layer with random phase sources above a metal ground plane. The directivity of emission from the random sources is lost.

consider the band structure in the $k_x$–$k_y$ plane. This means that at 54.3 THz, we can obtain a directional emission beam not just in the $H$-plane but also in all directions, although the structure is periodic in only one direction.

We note that there are other schemes to realize coherent thermal radiations, such as by utilizing three-dimensional (3D) photonic crystals [36] or 1D photonic crystal cavities [37]. Our metamaterial slab presents a route to achieve linearly polarized coherent thermal emission radiations in a wide frequency range, which can be tuned by adjusting structural parameters and material parameters.
5. Conclusion

In summary, we have proposed a simple metamaterial slab structure that possesses spatially coherent magnetic surface resonance states in a broad range of frequencies. These states facilitate nearly perfect absorption in a thin metamaterial slab containing slightly absorptive materials. As the absorption spectrum is highly angle-selective, the slab should give directional thermal emission. Direct FDTD simulation with random-phase sources corroborates the existence of strong angular emissivity even in the presence of structural disorder. As the surface resonances originate from artificial resonators, the operational frequency and the response can be tuned by varying the structural configurations. The simple metamaterial structure may be a useful platform to realize the coherent control of thermal emissions, optical antennas, IR or THz spectroscopy, photon detectors and active plasmonic devices.

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