Plasmon polaritons in photonic metamaterial superlattices: Absorption effects

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We discuss the propagation of electromagnetic waves in layered structures made up of alternate layers of air and metamaterials. The role played by absorption on the existence of electric and magnetic plasmon polaritons is investigated. Results show that plasmon-polariton modes are robust even in the presence of rather large absorption.

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The physical realization of metamaterial photonic systems [1,2] stimulated worldwide activity on wave propagation [3–5] in left-handed materials (LHMs). It opened up new possibilities to manipulate light, as metamaterial heterostructures exhibit unusual properties such as reverse Doppler shift, reverse Cherenkov radiation, reverse refraction, superlenses that amplify evanescent waves and overcome the diffraction limit [6], and bandpass optical filters with huge roll factors [7].

Simple one-dimensional (1D) layered structures containing LHMs also exhibit novel features, such as enhanced gaps, non-Bragg gaps corresponding to zero average refractive index [8], the possibility of a complete three-dimensional (3D) band gap in a 1D periodic structure [9], and unusual Anderson localization properties [10–12]. LHMs have not only opened a new era for optical devices, but have also given considerable thrust to a new branch of photonics termed *plasmonics* which deals with the resonant coupling between light and matter in nanostructured superlattices [13].

Recently, a study on the oblique incidence of light in 1D stacks composed of nondissipative alternating layers of air and LHMs has reported the existence of plasmon-polariton modes. They are of electrical nature for transverse magnetic (TM) waves or of magnetic nature for transverse electric (TE) waves [14], and appear in the spectral region where both dielectric-permittivity and magnetic-permeability responses of the dispersive metamaterial are zero. In the former case, they are coupled modes that result from resonant interactions between electromagnetic waves and mobile electrons of the metamaterial. In the latter, they correspond to resonant interactions between electromagnetic waves and magnetic current densities within the metamaterial.

This Brief Report considers the problem of existence of plasmon polaritons in lossy 1D structures. It is known that left-handed photonic systems are intrinsically absorptive. Also, one has all reasons to believe that the existence of plasmon-polariton modes could be hindered by absorption. Indeed, in the lossless case they are expected to appear near the band edge, i.e., precisely in the frequency regions where bands are most strongly affected by introduction of losses. Here we demonstrate that, somewhat counterintuitively, even quite strong absorption does not prevent an existence of plasmon polaritons.

Let us consider the oblique incidence of light on a 1D superlattice composed of layers A of air, and layers B of a doubly negative, dispersive, and absorptive material. Layers A (dielectric permittivity, $\epsilon_a=1$, magnetic permeability, $\mu_a=1$, and width a) and B (dispersive absorptive LHM, and width b) are distributed periodically so that $d=a+b$ is the period of the superlattice nanostructure.

We choose a Drude-like response for both the dielectric permittivity and magnetic permeability of the metamaterial component of the stack [3,4],

$$\epsilon_b(\omega)=\epsilon_0-\frac{\omega_e^2}{\omega(\omega+i\gamma_e)},\quad \mu_b(\omega)=\mu_0-\frac{\omega_m^2}{\omega(\omega+i\gamma_m)},\tag{1}$$

where $\nu_e=\frac{\omega_e}{2\pi\sqrt{\epsilon_0}}$ and $\nu_m=\frac{\omega_m}{2\pi\sqrt{\mu_0}}$ are the frequencies associated to the electric and magnetic plasmon modes, respectively, and $\gamma_{e,m}$ are damping constants.

From the transfer-matrix method, one obtains the transmission coefficient, T, and the reflection coefficient, R, for a finite stack composed of N double-stack periods, the following expressions:

$$T=\frac{2\xi}{\xi(M_{22}+M_{11})-\xi^2M_{12}-M_{21}},\tag{2}$$

$$R=\frac{\xi(M_{11}-M_{22})-\xi^2M_{12}+M_{21}}{\xi(M_{22}+M_{11})-\xi^2M_{12}-M_{21}},\tag{3}$$

where $\xi=\cos\theta$, and θ is the incidence angle. The M_{ij} are elements of the transfer matrix,

$$\mathbf{M}=\begin{bmatrix}M_{11} & M_{12}\\ M_{21} & M_{22}\end{bmatrix},\ \mathbf{M}=\prod_{j=1}^{N}\mathbf{M}_j^{(a)}\mathbf{M}_j^{(b)}.\tag{4}$$

Here, matrices $\mathbf{M}_j^{(a)}$ and $\mathbf{M}_j^{(b)}$ correspond to air (a) and LHM (b) layers, respectively, of the jth double-stack cell. One has

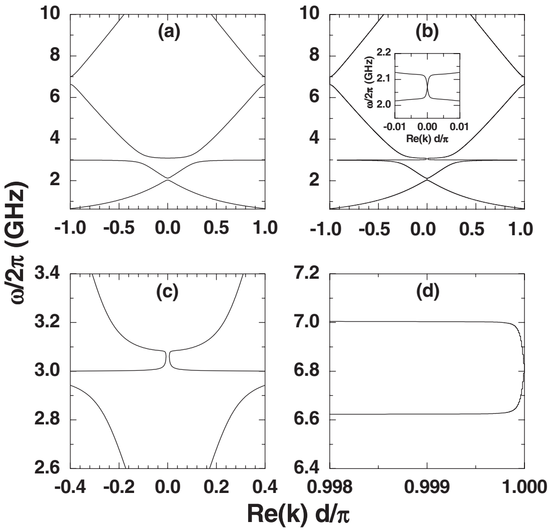

FIG. 1. TE dispersion $\nu=\nu(k)$ corresponding to incident light at an angle $\theta=\pi/12$ in a photonic periodic superlattice ($\nu=\frac{\omega}{2\pi}$). We have followed previous work [3] and chosen $\epsilon_0=1.21$ and $\mu_0=1.0$ in Eq. (1). Calculations were performed for air as slab A ($\epsilon_A=1$, $\mu_A=1$), $a=b=12$ mm, and $\omega_e/2\pi=\omega_m/2\pi=3$ GHz for the Drude model in slab B. In (a), an ideal nondissipative system is presented, i.e., $\gamma_{e,m}=0$, whereas in (b), (c), and (d) absorption is included with $\frac{\gamma_{e,m}}{\omega_{e,m}}=0.001$. Note the inset in (b), as well as the amplifications of (b) shown in (c) and (d).

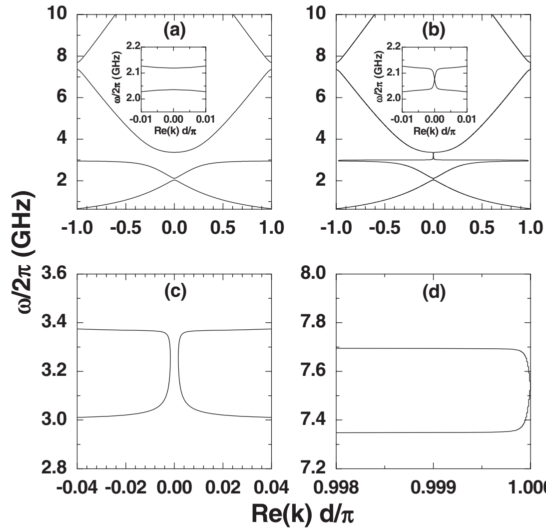

FIG. 2. Same as in Fig. 1, with incidence angle $\theta=\frac{\pi}{6}$.

$$\mathbf{M}_j^{(\alpha)}=\begin{bmatrix}\cos(q_\alpha\alpha) & i\sin(q_\alpha\alpha)/f_\alpha\\ if_\alpha\sin(q_\alpha\alpha) & \cos(q_\alpha\alpha)\end{bmatrix}, \tag{5}$$

where $\alpha=a,b$ and $q_\alpha=\omega/c\sqrt{\epsilon_\alpha(\omega)\mu_\alpha(\omega)-\sin^2\theta}$. For incident TE and TM waves, the coefficients f_α are given by

$$f_\alpha^{\mathrm{TE}}=\frac{1}{\mu_\alpha(\omega)}\sqrt{\epsilon_\alpha(\omega)\mu_\alpha(\omega)-\sin^2\theta},$$

$$f_\alpha^{\mathrm{TM}}=\frac{1}{\epsilon_\alpha(\omega)}\sqrt{\epsilon_\alpha(\omega)\mu_\alpha(\omega)-\sin^2\theta}. \tag{6}$$

For the case of an infinite periodic structure, Eqs. (5) and (6) lead to the following dispersion relation:

$$\cos(kd)=\cos(q_a a)\cos(q_b b)-\frac{Z}{2}\sin(q_a a)\sin(q_b b), \tag{7}$$

where k denotes the component of the wave vector in the growth direction and $Z=f_a/f_b+f_b/f_a$.

For the sake of comparison, let us begin by discussing the nondissipative case, that is, the one with real responses [$\gamma_e=\gamma_m=0$ in Eq. (1)]. Here, we concentrate on the analysis of TE polarization, but similar results hold for the TM case. Equation (7) gives a clear picture of the spectra of an infinite periodic stack. Figures 1(a) and 2(a) show three different band gaps, i.e., the non-Bragg $\langle n\rangle=0$ gap corresponding to the zero average refractive index defined as

$$\langle n\rangle=\frac{an_a+bn_b}{a+b},\quad n_\alpha=\sqrt{\epsilon_\alpha\mu_\alpha}$$

(the lower one), the plasmon-polariton gap around the frequency region where both dielectric permittivity and magnetic permeability are zero, and the usual Bragg gap at the Brillouin zone boundary (the upper one). Properties of these gaps are rather different. Since changes in the incidence angle affect strongly both the polariton and Bragg gaps, we illustrated our consideration for two different angles of incidence ($\theta=\pi/12$ for Fig. 1 and $\theta=\pi/6$ for Fig. 2).

In the nondissipative case [Figs. 1(a) and 2(a)], the non-Bragg $\langle n\rangle=0$ gap persists for all angles of incidence as clearly shown in previous work [14]. It becomes narrower with increasing of the incidence angle. In the TE case, the plasmon-polariton gap, for oblique incidence, corresponds to plasmon polaritons of a magnetic nature. Near the edges of the Brillouin zone, this gap tends to be of a plasmonlike nature, for the lower branch, while of a photonlike nature, for the upper branch. On the other hand, in the neighborhood of the center of the Brillouin zone the interpretation is reversed, i.e., the lower band corresponds to a photonlike mode and the upper band is a plasmon-type excitation. Contrary to the non-Bragg $\langle n\rangle=0$ gap, polariton gaps become wider with increasing of the incidence angle.

Let us now turn to case when absorption is present [one can see it in Figs. 1 and 2(b)–2(d)]. It is well known that presence of arbitrarily small absorption makes an existence of a complete band gap impossible. One can see that for small losses, while the plasmon-polariton modes persist essentially as in the nondissipative case around the border of the Brillouin zone, the gap disappears at the center of the Brillouin zone [cf. Figs. 1(b) and 2(b) and amplifications in Figs. 1(c) and 2(c)]. Also, both non-Bragg and Bragg gaps disappear with absorption, as seen in the insets in Figs. 1(b) and 2(b), and Figs. 1(d) and 2(d), respectively. It is interesting that for larger incidence angles deformations of bands in the vicinity of frequency regions corresponding to polaritonic gaps of lossless structures are smaller for both edges

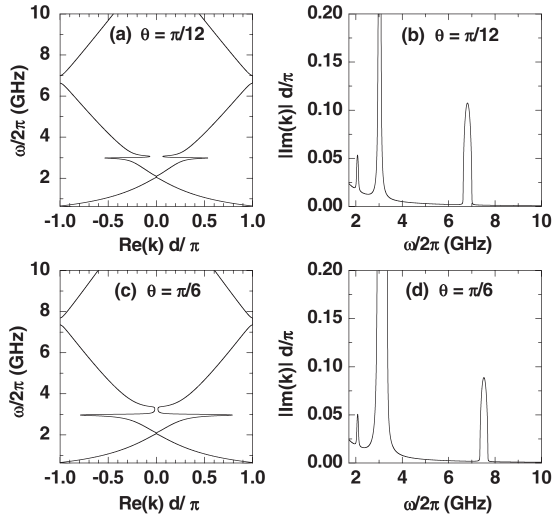

FIG. 3. TE dispersion relations [(a) and (c)], and imaginary parts of the wave vector [(b) and (d)], for incidence angles $\theta = \pi/12$ [(a) and (b)] and $\theta = \pi/6$ [(c) and (d)]. Here $\frac{\gamma_{e,m}}{\omega_{e,m}} = 0.01$. Parameters as in Fig. 1.

and centers of the Brillouin zone than for smaller incidence angles (see Fig. 3 and 4). However, even for large absorption, when bands are deformed to the point of nearly complete obliteration of features resembling a band-gap for lossless structure (for example, Fig. 4 corresponds to $\frac{\gamma_{e,m}}{\omega_{e,m}} = 0.1$) plasmon-polariton modes can still exist and influence significantly wave propagation through the structure. It becomes obvious when looking on imaginary part of solutions for the wave-vector given by Eq. (7). One can understand it by considering the lossless case. For example, for TE waves and magnetic plasmon polaritons in the lossless case from the dispersion relation (7) it follows that

$$\mathrm{Im}(k) \sim -\{\mathrm{Im}[\mu_b(\omega)]\}^{-1} \to -\infty,$$

when

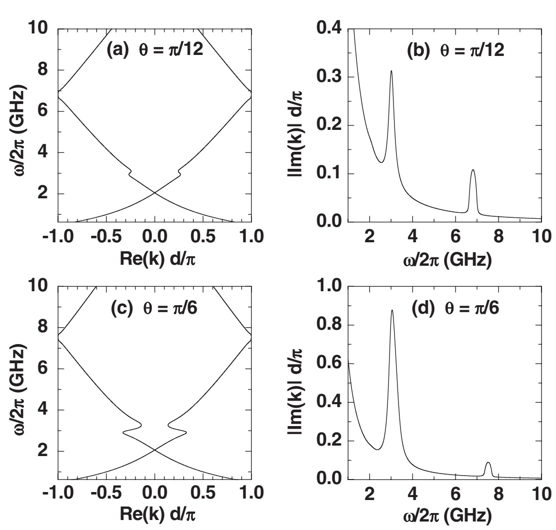

FIG. 4. Same as in Fig. 4, with absorption given by $\frac{\gamma_{e,m}}{\omega_{e,m}} = 0.1$

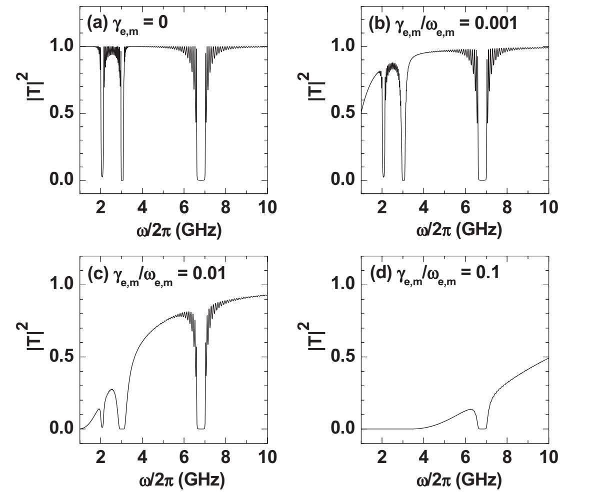

FIG. 5. Transmission for a photonic system containing $N=50$ double-stack layers in a TE configuration with incidence angle $\theta = \pi/12$, without (a) and with absorption for $\frac{\gamma_{e,m}}{\omega_{e,m}} = 0.001$ (b), $\frac{\gamma_{e,m}}{\omega_{e,m}} = 0.01$ (c) and $\frac{\gamma_{e,m}}{\omega_{e,m}} = 0.1$. Other parameters as in Fig. 1.

$$\mathrm{Re}[\mu_b(\omega)] = 0 \quad \text{and} \quad \mathrm{Im}[\mu_b(\omega)] \to +0.$$

Due to infinite contrast between refractive indices of LHM and air a total internal reflection will take place for an arbitrary small (but finite) incidence angles, and the field penetration beyond the LHM-air interface would be very small. So, in the case of sufficiently small absorption present one can still expect rather weak penetration of the incident field to the structure and, consequently, low absorption and high reflection of the field. Rather large values of imaginary parts of the Bloch wave-vector in region of the polariton gaps of lossless structures are present even for significant absorption (see Fig. 4). It indicates that in this case the picture outlined above will hold, too.

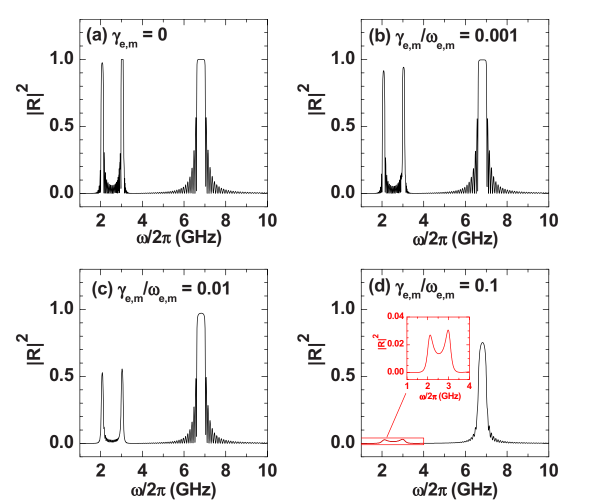

FIG. 6. (Color online) Reflection for a photonic system containing $N=50$ double-stack layers in a TE configuration with the incidence angle $\theta = \pi/12$, without (a) and with absorption for $\frac{\gamma_{e,m}}{\omega_{e,m}} = 0.001$ (b), $\frac{\gamma_{e,m}}{\omega_{e,m}} = 0.01$ (c) and $\frac{\gamma_e}{\gamma_m} = 0.1$. Other parameters as in Fig. 1.

Transmission/reflection studies confirm that prognosis. We consider the transmission and reflection spectra of a plane wave incident on a finite 1D stack made up of $N=50$ double-layer cells. Transmission and reflection results are plotted in Figs. 5 and 6, for an incidence angle of $\frac{\pi}{12}$, and for $\gamma_{e,m}=0$, $\frac{\gamma_{e,m}}{\omega_{e,m}}=0.001$, $\frac{\gamma_{e,m}}{\omega_{e,m}}=0.01$, and $\frac{\gamma_{e,m}}{\omega_{e,m}}=0.1$. As expected, in the nondissipative case, the patterns of transmission/reflection clearly follow the band-gap structures exhibited in Figs. 1–4. When damping is taken into account, we note that transmission and reflection features at the frequencies corresponding to the non-Bragg, plasmon-polariton, and Bragg gaps are broadened and smoothed out. We note that even for $\frac{\gamma_{e,m}}{\omega_{e,m}}=0.1$, the small reflection intensities depicted in the inset of Fig. 6(d) are still large enough to be experimentally observed. We re-emphasize that, even for rather significant levels of absorption ($\frac{\gamma_{e,m}}{\omega_{e,m}}=0.1$), both the non-Bragg and the plasmon-polariton structures are distinctly observable.

To conclude, we have considered dispersion relations for a periodic 1D double-stack photonic superlattice containing a metamaterial layer. We explicitly take into account absorption effects for the metamaterial. We have verified that the coupled plasmon-polariton modes survive, and may be experimentally observed. Analogously to the plasmon-polariton of electric nature, that is able to confine light to very small dimensions, the existence of these plasmon-polariton modes of a magnetic nature could well be exploited to produce ultracompact photonic devices for sensing, guiding or focusing.

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