Perfect Absorption in Ultrathin Epsilon-Near-Zero Metamaterials Induced by Fast-Wave Non-Radiative Modes

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Above-light-line surface plasmon polaritons can arise at the interface between a metal and ε-near-zero metamaterial. This unique feature induces unusual fast-wave non-radiative modes in a ε-near-zero material/metal bilayer. Excitation of this peculiar mode leads to wide-angle perfect absorption in low-loss ultrathin metamaterials. The ratio of the perfect absorption wavelength to the thickness of the ε-near-zero metamaterial can be as high as 104; the electromagnetic energy can be confined in a layer as thin as λ/10000. Unlike conventional fast-wave leaky modes, these fast-wave non-radiative modes have quasi-static capacitive features that naturally match with the space-wave field, and thus are easily accessible from free space. The perfect absorption wavelength can be tuned from mid- to far-infrared by tuning the ε ≈ 0 wavelength while keeping the thickness of the structure unchanged.

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A metamaterial is a composite structure with an electromagnetic (EM) response not readily observed in naturally occurring materials. Many remarkable phenomena have been predicted [1–4] by tuning the permittivity ε and permeability μ in extraordinary ways. If the dielectric response is made vanishingly small, creating an epsilon-near-zero (ENZ) material, interesting radiative effects are expected to occur [5–7]. On the other hand, by manipulating the EM response to achieve a small transmittance (T) and reflectance (R), enhanced absorption can ensue. Strong absorption in a thin layer typically requires high loss. One of the earliest absorbers, the Salisbury screen [8], is based on the phenomenon of destructive wave interference, and is thus limited to a minimum thickness of one quarter wavelength. To overcome the thickness constraint, absorbing screens using metamaterials [9] and high impedance ground planes [10] have recently been proposed. By exploring resonant enhancement, thin metamaterial and nanoplasmonic absorbers were demonstrated in structures having localized resonances [11–23]. In those structures, the geometrical quality factor (GQF), i.e. the ratio of the perfect absorption wavelength to the thickness of the medium, is significantly improved compared to the standard Salisbury screen. The best GQF of those structures is about 40 [14].

Based on a fundamentally different mechanism, in this paper we demonstrate wide-angle perfect absorption in a low-loss ultrathin ENZ-metamaterial/metal bilayer as shown in Fig. 1. In our structure, the GQF can be as high as 104. This structure possesses fast-wave non-radiative (FWNR) modes due to unconventional above-light-line surface plasmon polaritons (ALL-SPPs) at the ENZ-metal interface, which is easily accessible from free space. Fast waves have phase velocities exceeding the speed of light in vacuum. Conventional fast waves in planar structures are radiative leaky modes that cannot be excited by plane wave incidence due to wave-vector mismatch at the boundary. For our ENZ-metal structure, the exotic FWNR mode is naturally matched to free space at the ENZ-air interface, and thus is easily accessible from free space. Unlike the Salisbury screen where the thickness of the absorbers increases with the absorption wavelength, in our structure the wavelength can be tuned from mid- to far-infrared (IR) by tuning the ε ≈ 0 wavelength while maintaining the thickness.

To investigate this absorption phenomenon in a general fashion, an anisotropic ENZ metamaterial (ε1x = 1 and ε1z ≈ 0) ENZ/metal bilayer with thicknesses d1 and d2 for the ENZ medium and metal (Ag), respectively. Above-light-line surface plasmon polaritons (ALL-SPPs) are excited at the ENZ-metal interface.

FIG. 1. (Color online) A plane wave is incident (at an angle θ) on an anisotropic (ε1x = 1 and ε1z ≈ 0) ENZ/metal bilayer with thicknesses d1 and d2 for the ENZ medium and metal (Ag), respectively. Above-light-line surface plasmon polaritons (ALL-SPPs) are excited at the ENZ-metal interface.

\[
\frac{k_{\text{SPP}}}{k_0} = \sqrt{\frac{\varepsilon_1\varepsilon_2\varepsilon_{2z}\varepsilon_{1x} - \varepsilon_1\mu_1\mu_2}{\varepsilon_1\varepsilon_2\varepsilon_{2z} - \varepsilon_1\varepsilon_1\varepsilon_2}} \approx \sqrt{\frac{\varepsilon_1\mu_1(1 - \delta)}{\varepsilon_1\mu_1}}, \quad (1)
\]

where \(k_0 = \omega/c\) and \(\delta = \varepsilon_1\mu_2/(\varepsilon_2\mu_1)\). The approximation is taken when \(\varepsilon_1 \to 0\). In our case, \(\mu_1 = \mu_2 = 1\) and \(\varepsilon_2 = \varepsilon_{2z} = \varepsilon_m \sim -10^3\) (in the IR region). Thus, \(k_{\text{SPP}}/k_0 \approx \sqrt{\varepsilon_1\mu_1} < 1\), which characterizes above light-line SPP dispersion, and is clearly different from conventional SPP dispersion.
The absorption can be calculated from Maxwell’s equations. Assuming a harmonic time dependence $\exp(-i\omega t)$ for the EM field, we have
\[
\nabla \times (\tilde{\mu}_n^{-1} \cdot \nabla \times E) = \kappa_0^2 (\tilde{\varepsilon}_n \cdot E),
\]
\[
\nabla \times (\tilde{\varepsilon}_n^{-1} \cdot \nabla \times H) = \kappa_0^2 (\tilde{\mu}_n \cdot H),
\]
where $\tilde{\varepsilon}_n$ and $\tilde{\mu}_n$ are, respectively, the permittivity and permeability tensors for a given (uniform) region ($n = 0, 1, 2, \cdots$), which in the principal coordinates are described by,
\[
\tilde{\varepsilon}_n = \varepsilon_{nx} \hat{x} \hat{x} + \varepsilon_{ny} \hat{y} \hat{y} + \varepsilon_{nz} \hat{z} \hat{z},
\]
\[
\tilde{\mu}_n = \mu_{nx} \hat{x} \hat{x} + \mu_{ny} \hat{y} \hat{y} + \mu_{nz} \hat{z} \hat{z}.
\]

We consider TM modes, corresponding to non-zero field components $H_y$, $E_x$, and $E_z$. The magnetic field $H_y$ satisfies the following wave equation:
\[
\frac{1}{\varepsilon_x} \frac{\partial^2 H_y}{\partial x^2} + \frac{1}{\varepsilon_x} \frac{\partial^2 H_y}{\partial z^2} + k_0^2 \mu_y H_y = 0,
\]
which admits solutions of the form $\psi(z) \exp(i\beta x)$. The parallel wave-vector $\beta$ is determined by the incident wave, and is conserved across the interface,
\[
\beta^2 = k_0^2 \varepsilon_{nx} \mu_y - \alpha_n^2 \varepsilon_{nx}, \quad (n = 0, 1, 2, \cdots),
\]
where $\alpha_n$ is the wave number in the $z$ direction. The functional form of $\psi(z)$ is either a simple exponential $\exp(i\alpha_n z)$ for the semi-infinite region or a superposition of $\cos(\alpha_n z)$ and $\sin(\alpha_n z)$ terms for the bounded region (along the $z$ direction). The other two components $E_x$ and $E_z$ are found using Maxwell’s equations. By matching boundary conditions at the interface, i.e., continuity of $H_y$ and $E_z$, the transmittance ($T$) and reflectance ($R$) can be calculated via the Poynting vector $S$, given by $S = c/8\pi \mathbf{R} \cdot (E \times H^*)$. The fraction of energy that is absorbed by the system is determined by the absorptance ($A$), where $A = 1 - T - R$, consistent with energy conservation. Figure 2 shows the absorptance of the ENZ-metal structure versus the angle of incidence (AOI) in the mid- to far-IR regime. Unless stated otherwise, the results that follow correspond to $\varepsilon_{nx} = 0.001 + i 0.01$. The AOI at which the perfect absorption occurs depends on the wavelength and the thickness of the ENZ layer. The full-width-half-maximum (FWHM) angular width can be as high as $\sim 60^\circ$ (top-left panel), while the GQF can reach $\sim 10^\circ$ (top-right panel). When the thickness increases, the angular bandwidth varies differently in the mid- and far-IR regions.

In Fig. 3, we show how the absorptance varies with the AOI and incident wavelength (right panels). The left set of panels corresponds to the incident angle $\theta_p$ that results in perfect absorption, as a function of wavelength. Perfect absorption occurs when the $\theta$ dependent effective impedance ($Z_e$) of the ENZ-metal structure matches that of free space ($Z_0$), i.e. $Z_e = Z_0$, in which case the reflection coefficient is zero. The effective impedance of the ENZ-metal structure can be expressed as
\[
Z_e = Z_1 \frac{1 - r_{12} \exp(i2\phi)}{1 + r_{12} \exp(i2\phi)} = Z_0,
\]
\[
\tan \phi = -\frac{Z_1 (Z_0 + Z_2)}{Z_1 + Z_0 Z_2},
\]
which implicitly depends on the parallel complex propagation constant $\beta$ (Eq. (5)). The mode solutions of Eq. (7) have four branches due to the $\pm$ signs inherent to the square root of $\alpha_0$ and $\alpha_2$ in Eq. (5). The branch leading to perfect absorption corresponds to both $\alpha_0 < 0$ and $\alpha_2 < 0$. In the long wavelength regime and for ultrathin slabs, this branch provides fast-wave ($\beta/k_0 < 1$), non-radiative ($\beta$ real) modes that
are guided along the ENZ layer. In the left panels of Fig. 3 we compute the direction of perfect absorption as a function of wavelength in three complimentary ways: The solid (blue) curves are numerically extracted from the absorptance corresponding to the right panels, the dashed (green) curves are obtained from Eq. (6), where the structure is impedance matched to free space, and the circles (red) are calculated using Eq. (7) and a root searching algorithm. The overlap of each set of data further demonstrates that the underlying mechanism of perfect absorption is the excitation of FWNR modes. We see also that perfect absorption occurs only at oblique incidence where the electric field has a nonzero component in the normal direction.

The energy density, $U$, in lossy and dispersive anisotropic media is given by,

$$U = \frac{1}{16\pi} \Re \left[ E^\dagger \cdot \frac{\partial (\omega E)}{\partial \omega} E + H^\dagger \cdot \frac{\partial (\omega H)}{\partial \omega} H \right]. \quad (8)$$

We show in Fig. 4 the spatial distribution of the energy density in the ENZ-metal structure under conditions of perfect absorption at mid- and far-IR wavelengths.

**FIG. 4.** (Color online) Spatial distribution of the energy density computed from Eq. (8) (normalized so that $H_z$ is unity at the ENZ-air interface) for $\lambda = 10 \mu m$ and AOI = 16.1° (left) and $\lambda = 155 \mu m$ and AOI = 56.3° (right). The ENZ-layer with $d_l = 0.2 \mu m$ is located at $z \in [0.1, 0.1]$. The regions $z \in [0.1, 0.15]$ and $z \in [-0.2, -0.1]$ are, respectively, silver and air. Color-bars represent the magnitude ($\times 0.01$) of the energy density, indicating the energy is more confined to the ENZ medium at far-IR wavelengths.

For the minimum thickness of ENZ materials is generally limited by fabrication methods except in some cases of naturally occurring resonances. Figure 5 shows the direction, $\beta_p$ (top left), and FWNR angular bandwidth (bottom left) of perfect absorption vs. ENZ thickness. The curves are numerically extracted from the absorptance when $\lambda = 10 \mu m$ (dots (blue)), $\lambda = 100 \mu m$ (solid (green)), and $\lambda = 300 \mu m$ (dashed (red)). The circles in the left-top panel are obtained from Eq. (7). The normalized propagation constant $\beta/k_0$ corresponding to these modes is shown in the right panels. Top-right: real part of $\beta/k_0$. Bottom-right: imaginary part of $\beta/k_0$. The Ag thickness ($d_L$) is 10 nm.

**FIG. 5.** (Color online) Direction (top left) and FWHM bandwidth (bottom left) of perfect absorption vs. ENZ thickness. The curves are numerically extracted from the absorptance when $\lambda = 10 \mu m$ (dots (blue)), $\lambda = 100 \mu m$ (solid (green)), and $\lambda = 300 \mu m$ (dashed (red)). The circles in the left-top panel are obtained from Eq. (7). The normalized propagation constant $\beta/k_0$ corresponding to these modes is shown in the right panels. Top-right: real part of $\beta/k_0$. Bottom-right: imaginary part of $\beta/k_0$. The Ag thickness ($d_L$) is 10 nm.

The average energy density in the ENZ medium is also calculated (Eq. (8)), as well as the effective surface charge density $Q = \epsilon_{\parallel} E_z/4\pi$ (normalized by the input $E_z$) and capacitance $C = Q/V$ per unit area, which is much larger than the capacitance ($C_0 = \epsilon_{\parallel}/d_l$) of a parallel-plate capacitor (of unit-area), leading to perfect absorption in the ultrathin ENZ medium.
for different where shows the ratio of transport power to dissipation power, (blue), when constant behavior of the charge (right-top) as a function of panels). The linear responses of the voltage (middle-top) and panels). The left and right set of panels correspond to material. These results are presented in Fig. shows the ratio of transport power to dissipation power, suggesting longer wavelengths have better power handling capability if the structure is used as an ultrathin channel to transport light. The influence of loss on the absorption is illustrated in Fig. 7 for different $\epsilon_1$ (real part) in the mid- and far-IR. In the far-IR, the smaller $\Im(\varepsilon_1)$ has larger angular bandwidth, which is opposite to what occurs in the mid-IR.

In summary, wide-angle perfect absorbers have been demonstrated in ultrathin ENZ-metal structures, and which have extraordinary high geometric quality factors. Such structures may have applications involving ultra light-weight infrared absorbers, cloaking, EM shielding, photovoltaic systems, and highly efficient infrared sensors and detectors. The authors gratefully acknowledge the sponsorship of ONR N-STAR and NAVAIR’s ILIR programs.

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