Coherent Control of Low Loss Surface Polaritons

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We propose fast all-optical control of surface polaritons by placing an electromagnetically induced transparency (EIT) medium at an interface between two materials. EIT provides longitudinal compression and a slow group velocity, while matching properties of the two materials at the interface provides strong transverse confinement. In particular, we show that an EIT medium near the interface between a dielectric and a negative-index metamaterial can establish tight longitudinal and transverse confinement plus extreme slowing of surface polaritons, in both transverse electric and transverse magnetic polarizations, while simultaneously avoiding losses.

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Introduction.—All-optical rapid guidance, processing, and control of light in nanophotonic and quantum information applications is important but limited by weak nonlinearities in typical materials, the fast speed of light, and physical limitations to confinement such as enhanced absorption and the narrow spectral width of light. Strategies to overcome these speed and confinement weaknesses include exploiting photonic crystals with defect structures to trap light [1], surface polaritons (SPs) to confine light to wavelength dimensions [2] and use nonlinear interactions [3,4], and electromagnetically induced transparency (EIT) to slow and compress light in the longitudinal propagation direction [5], which are realized in the solid state [6], Bose-Einstein condensates [7], and Mott insulators [8]. We propose placing an EIT medium at the interface of two materials. This arrangement benefits from the combination of transverse confinement of surface polaritons with the longitudinal compression and slowing of pulses due to EIT.

We derive an analytical solution to the problem, which provides an elegant picture of EIT at a material interface, including hitherto unsuspected properties for the interface between a dielectric and a negative-index metamaterial (NIMM) [9]: Our two key results for the dielectric-NIMM interface are (i) that this interface simultaneously supports both transverse electric (TE) and transverse magnetic (TM) polarizations, whereas dielectric-dielectric and dielectric-metal interfaces support only TM, and (ii) SP loss can be made arbitrarily small for the dielectric-NIMM interface but not for the other cases. Our analysis applies to general interfaces, but the dielectric-NIMM interface is especially intriguing and suggests low loss, fast control of both TM and TE SPs.

For our scheme in Fig. 1, we suggest end-fire coupling excitation, which has been demonstrated experimentally with a high efficiency of 0.7 [10]. Our approach is to create two SP fields by directing two laser beams at the interface between two materials with permittivities ε_i and permeabilities μ_i (top: $i = 1$; bottom: $i = 2$), and these two SP fields propagate toward the EIT medium comprising three-level Λ atoms (3LA), quantum dots, nitrogen-valence centers in diamond, rare-earth ions in crystals, or similar (henceforth the "Λ medium"); the 3 levels are designated $|\ell\rangle$ for $\ell \in \{1, 2, 3\}$, and the transition frequency $\omega_{\ell\ell'}$ corresponds to $|\ell\rangle \leftrightarrow |\ell'\rangle$. Our analysis does not restrict the signs of ε and μ and hence accommodates dielectrics, surface plasmons at a dielectric-metal interface, and also NIMMs where both $\varepsilon, \mu < 0$.

Absorption and dispersion of SP fields.—The SP fields, which propagate in the positive x direction, are characterized by the ω-dependent complex wave vector $k_\parallel + i\kappa$. For $k_j^2 = k_\parallel^2 - \omega^2\varepsilon_j\mu_j/c^2$, the wave vector component normal to the interface, $k_1\varsigma_2 = -k_2\varsigma_1$, where $\varsigma \equiv \mu$ for the TE

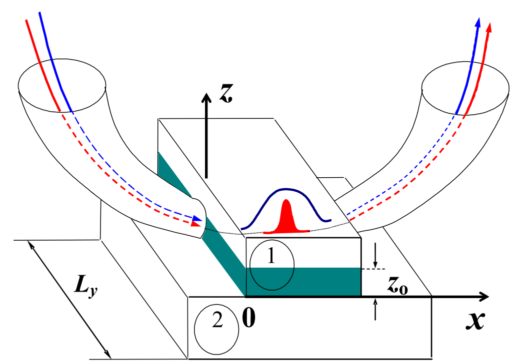

FIG. 1 (color online). End-fire coupling scheme for coherent control of SPs. Two incoming beams from the left waveguide create two SP pulses above the interface between media 1 and 2 in half-spaces $z > 0$ and $z < 0$, respectively: control (blue line) and signal probe (red area). The SPs interact with a collection of Λ medium shown above the interface as a shaded green layer of thickness z_0.

mode and $\varsigma \equiv \varepsilon$ for the TM mode; both polarizations coexist only if both conditions are simultaneously satisfied.

We analyze the SP modes at the interface between a dielectric and a NIMM medium. Although the dielectric-NIMM interface is technically challenging, rapid progress is bringing metamaterials to the optical domain [11], thereby opening possibilities for optical storage and control [12]. The dielectric-NIMM interface is especially attractive because both TE and TM polarization modes can coexist, and, as we show below, complete suppression of SP losses is possible, in principle, thereby admitting novel opportunities in SP-field control. Optical properties of the NIMM can be modeled with complex dielectric permittivity and magnetic permeability given by [9,10] $\varsigma_2(\omega) = \varsigma_r + i\varsigma_i = 1 - \frac{\omega_f^2}{\omega(\omega + i\gamma_f)}$ for the two cases $\varsigma \equiv \varepsilon$, $f \equiv e$ and $\varsigma \equiv \mu$, $f \equiv m$. Here $\omega_{e,m}$ are the electric and magnetic plasma frequencies of the NIMM and $\gamma_{e,m}$ the loss rates, respectively. Accounting for complex ε and μ in the surface boundary conditions of SP wave vectors, we find (TM case)

$$k_{\parallel}(\omega) + i\kappa(\omega) = \frac{\omega}{c}\sqrt{\varepsilon_1\varepsilon_2(\varepsilon_2\mu_1 - \epsilon_1\mu_2)/(\varepsilon_2^2 - \varepsilon_1^2)}, \tag{1}$$

and the TE case is similar. The real (imaginary) part of Eq. (1) yields SP TM mode dispersion (loss).

We study the system numerically at room temperature and at optical frequencies to determine its features. For metal, we use the values for Ag [13]: $\omega_e = 1.37 \times 10^{16}\ \mathrm{s}^{-1}$, $\gamma_e = 2.73 \times 10^{13}\ \mathrm{s}^{-1}$, and $\varepsilon_1 = 1.3$, $\mu_1 = 1$. As NIMM technology is embryonic, we consider a wide range of magnetic plasmon frequency such that $\omega_m \leq 0.5\omega_e$ [10], and γ_m is between $10^{-5}\gamma_e$ and γ_e itself.

Numerical results for $\kappa(\omega)$ reveal a deep abyss frequency ω_0 with complete cancellation of loss $\kappa(\omega_0) \sim 0$ corresponding to a specific ratio of magnetic and electric loss for each ω. From Eq. (1), ω_0 is accurately determined from

$$\frac{\mu_i}{\varepsilon_i} = \frac{\mu_r(\varepsilon_r^2 + \varepsilon_1^2) - 2\varepsilon_r\varepsilon_1}{\varepsilon_r(\varepsilon_r^2 - \varepsilon_1^2)}, \tag{2}$$

which fails for a dielectric-metal interface because $\mu_1, \mu_2 > 0$. Figure 2 compares SP losses at a dielectric-NIMM interface for $\gamma_m = 10^{11}\ \mathrm{s}^{-1}$ and $\omega_m = 0.5\omega_p$ and surface plasmon polaritons at a dielectric-metal interface.

We observe that NIMMs enable absorption losses to be drastically reduced in a narrow frequency band, as studied for freely propagating fields [14]. Here we predict a similar effect but for SP fields using a typical NIMM interface. Equation (2) shows that $\varepsilon_r, \mu_r < 0$ implies destructive interference between electric and magnetic absorption responses, which explains the high suppression of losses around ω_0. The relative weakness of decay rates makes the frequency ω_0 sensitive to electric and magnetic decoherence rates as is evident from Fig. 3. Henceforth, we

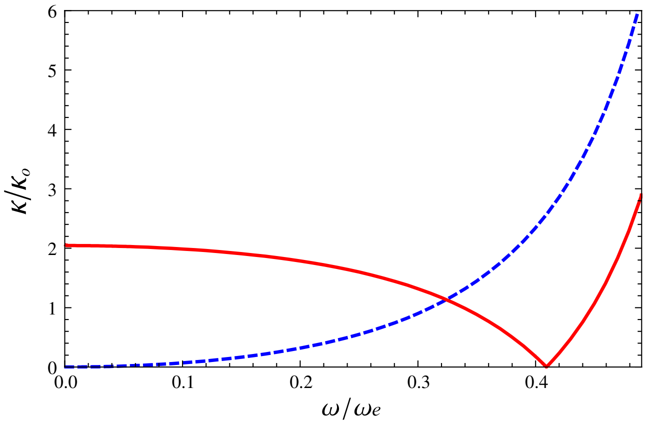

FIG. 2 (color online). Comparison of SP losses at a dielectric-NIMM interface (red curve) and losses of surface plasmons at a dielectric-metal interface (blue curve) for $\gamma_m = 10^{11}\ \mathrm{s}^{-1}$, $\gamma_e = 2.73 \times 10^{13}\ \mathrm{s}^{-1}$, $\omega_m = 0.5\omega_e$, $\omega_e = 1.37 \times 10^{16}\ \mathrm{s}^{-1}$, and $\omega_0 \cong 0.4092\omega_e$.

demonstrate the possibility of coherent control of slow SPs within frequency range around ω_0 of complete reduction of their losses in basic materials.

EIT control of SP modes.—Here we give a general analysis of EIT-based coherent control of SP pulse interacting with a Λ medium near the surface. We assume that the probe field in the TM mode has a frequency ω_{31} and control field frequency equal to ω_{32}. Transition $|1\rangle \leftrightarrow |2\rangle$ is dipole-forbidden.

The SP electric field near the surface is obtained from field quantization [15] in a dispersive medium [16]. In the plane wave expansion over modes λ,

$$\hat{\boldsymbol{E}}(x, z, t) = \sum_{\lambda} \int dk_{\parallel}[\boldsymbol{E}_{0\lambda}(k_{\parallel}, z)\hat{a}_{\lambda}(k_{\parallel}, t) \times e^{ik_{\parallel}x} + \mathrm{H.c.}],$$

$$[\hat{a}_{\lambda}(k_{\parallel}, t), \hat{a}^{\dagger}_{\lambda'}(k'_{\parallel}, t)] = 2\pi\delta_{\lambda\lambda'}\delta(k_{\parallel} - k'_{\parallel}).$$

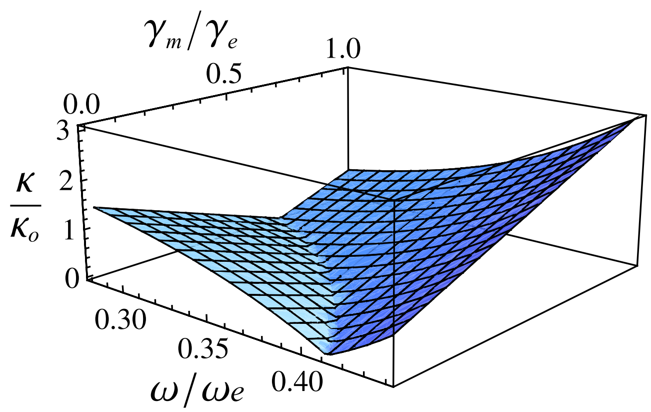

FIG. 3 (color online). Absorption loss for surface polaritons as a function of frequency ω/ω_e and magnetic decoherence rate γ_m/γ_e, where $\kappa_0 = 10^4\ \mathrm{m}^{-1}$, $\gamma_e = 2.73 \times 10^{13}\ \mathrm{s}^{-1}$, $\omega_m = 0.5\omega_e$, $\varepsilon_1 = 1.3$, and $\mu_1 = 1$.

We develop the theory for TM; the TE case is then a straightforward generalization. Dropping λ, we obtain

$$\boldsymbol{E}_0(k_\parallel, z) = \begin{cases} (\boldsymbol{e}_x + i\boldsymbol{e}_z k_\parallel/k_1)E_0(k_\parallel)e^{-k_1 z}, & z > 0, \\ (\boldsymbol{e}_x - i\boldsymbol{e}_z k_\parallel/k_2)E_0(k_\parallel)e^{k_2 z}, & z < 0. \end{cases} \quad (3)$$

Here $\boldsymbol{e}_x$ and $\boldsymbol{e}_z$ are unit vectors along the $\boldsymbol{x}$ and $\boldsymbol{z}$ directions, respectively, with electric field amplitude

$$E_0(k_\parallel) = \sqrt{\hbar\omega(k_\parallel)/2\pi\varepsilon_0 L_y L_z(\omega, \varepsilon, \mu)} \quad (4)$$

and transverse quantization length $L_z = D + \frac{\omega^2(k_\parallel)}{c^2}S$, with

$$D = \frac{\partial}{\partial\omega}(\omega\varepsilon_1)\frac{k_1^2 + k_\parallel^2}{k_1^3} + \frac{\partial}{\partial\omega}(\omega\varepsilon_2)\frac{k_2^2 + k_\parallel^2}{k_2^3},$$
$$S = \frac{\partial}{\partial\omega}(\omega\mu_1)\frac{\varepsilon_1^2}{k_1^3} + \frac{\partial}{\partial\omega}(\omega\mu_2)\frac{\varepsilon_2^2}{k_2^3}. \quad (5)$$

These quantities depend on the interface parameters and on the SP mode dispersion relation $\omega(k_\parallel)$ of Eq. (1).

Adopting the usual EIT approximations [4,17] for the evolution of a SP interacting with a Λ medium, we find the Fourier SP-field equation

$$(\partial/\partial x - i\nu/v_0)\hat{A}(\nu, x) = -[\alpha(\nu) + \kappa(\omega_{31})]\hat{A}(\nu, x) \quad (6)$$

on the surface: $\hat{A}(\nu, x) = (2\pi)^{-1}\int dt e^{i\nu t}\hat{A}(t, x, z = 0)$,

$$\hat{A}(t, x, z = 0) = e^{i\omega_{31}t - ik_\parallel^s(\omega_{31})x}\int dk_\parallel E_0(k_\parallel)\hat{a}(k_\parallel, t)e^{ik_\parallel^s x}.$$

Solving Eq. (6) yields the electric field over the surface

$$\hat{\boldsymbol{A}}(t, x, z > 0) = \left(\boldsymbol{e}_x + i\boldsymbol{e}_z\frac{k_\parallel(\omega_{31})}{k_1^s}\right)e^{-k_1^s z} \times \int d\nu e^{-i\nu t + [i(\nu/v_0) - \alpha(\nu) - \kappa(\omega_{31})]x}\hat{A}(\nu, 0), \quad (7)$$

where $\nu \equiv \omega(k_\parallel) - \omega_{31}$ is the SP-field detuning from the central frequency ω_{31}, which is assumed to be close to ω_0. Here $v_0 = \partial\omega/\partial k_\parallel$ is the SP group velocity without a 3LA Λ medium at $\omega = \omega_{31}$, and

$$\frac{\alpha(\nu)}{2\pi} = \frac{|g|^2}{v_0} \times \int_0^\infty \int_0^{L_y} dydz \frac{n(\boldsymbol{r})(\gamma_{21} - i\nu)e^{-2k_1^s z}}{|\Omega_c(\boldsymbol{r})|^2 - (\nu + i\gamma_{21})(\nu + i\Gamma_{31})} \quad (8)$$

yields dispersion and absorption of the SP field, with Γ_{31} a linewidth, $\Omega_c(\boldsymbol{r})$ the control field Rabi frequency, $g = \mathbf{d}\cdot(\boldsymbol{e}_x + i\boldsymbol{e}_z k_\parallel/k_1)E_0[k_\parallel(\omega_{31})]/\hbar$ the SP-Λ coupling constant, $\mathbf{d}$ the atomic dipole moment, and $n(\boldsymbol{r})$ the Λ medium density. These parameters can be optimized for SP-field control. Also, the control field, which yields Ω_c, can be a freely propagating mode or a SP TE or TM field. Henceforth, we assume $|\Omega_c(\boldsymbol{r})|^2 = |\Omega|^2 e^{-2k_1^c z}$, where k_1^s and k_1^c are the probe and control wave vectors, respectively, in the z direction for medium 1 as in Eq. (3).

Absorption and dispersion of the SP field in the presence of a 3LA medium is given by $\alpha(\nu) + \kappa(\omega_{31})$ with inhomogeneous broadening at resonance transition $h = \Delta_w/\pi(\Delta^2 + \Delta_w^2)$ and inhomogeneous broadening width Δ_w. Thus $\Gamma_{31} = \Delta_w + \gamma_{31}$, where we have assumed that all atoms share the same decay constants γ_{mn}; this assumption is reasonable for solid-state systems at liquid He temperatures because γ_{31}/Δ_w is negligible [4]. From Eq. (8), it follows that the field amplitude effectively is bounded by $z \leq \min\{1/k_1^s, 1/k_1^c\}$, and we assume confinement of the Λ medium to this height above the interface.

For ${}_2F_1$ the hypergeometric function Eq. (8) is integrable for constant density $n(\boldsymbol{r}) \equiv n(0 < z < z_0) = n$ yielding

$$\alpha(\nu, z_0) = \alpha_0(\omega_{31})G(k_1^s, k_1^c, z_0, \beta(\nu)), \quad (9)$$

$$G(k_1^s, k_1^c, z_0, \beta) = \frac{i\Gamma_{31}}{\nu + i\Gamma_{31}}\left[{}_2F_1\left(1, \frac{k_1^s}{k_1^c}, \frac{k_1^s + k_1^c}{k_1^c}, \frac{1}{\beta(\nu)}\right) - e^{-2k_1^s z_0}{}_2F_1\left(1, \frac{k_1^s}{k_1^c}, \frac{k_1^s + k_1^c}{k_1^c}, \frac{e^{-2k_1^c z_0}}{\beta(\nu)}\right)\right] \quad (10)$$

a spectral function, and $\beta \equiv (\nu + i\gamma_{21})(\nu + i\Gamma_{31})/|\Omega|^2$. Equation (10) is a maximum $G = 1$ for $\nu = 0$ and depends on Δ_w, k_1^s/k_1^c, z_0, and Ω_c, which provides rich opportunities for spectral and spatial control of the SP. For the thick atomic layer $z_0 \gg 1/k_1^s$, Eq. (10) simplifies: (i) $k_1^s/k_1^c \gg 1$ gives $G \approx i\beta(\nu)\Gamma_{31}/\{(\nu + i\Gamma_{31})[\beta(\nu) - 1]\}$ (typical EIT); (ii) $k_1^s/k_1^c = 1$ gives $G = -i\beta(\nu)\Gamma_{31}/(\nu + i\Gamma_{31}) \times \ln[1 - 1/\beta(\nu)]$ (modified EIT); and (iii) $k_1^s/k_1^c \ll 1$ gives $G \approx i\Gamma_{31}/(\nu + i\Gamma_{31})$ (no EIT).

As the formula is integrable, the resonant absorption coefficient for the Λ medium at ω_{31} is expressed in a simple form if the control field is off and $z_0 \gg 1/k_1^s$:

$$\alpha_0(\omega_{31}) = \pi n L_y |g|^2/k_1^s v_0(\omega_{31})\Gamma_{31}, \quad (11)$$

with coupling constant $|g|^2 \sim 1/L_z$. An appropriate choice of materials gives small L_z and hence considerably enhances the SP-field amplitude and increases interaction coupling between the SP field and the 3LA Λ medium.

Numerical example.—Using our solution we demonstrate the exciting possibility of EIT control for the *low loss SP modes*, where the low loss is due to interference between electric and magnetic responses, independent of EIT. For $L_y = 2.5\ \mu$m, $0.4087\omega_e < \omega_{31} < 0.4097\omega_e$, we find $\kappa(\omega_{31}) < 0.01\kappa_0$, $v_0(\omega_{31}) \approx 0.6c$, and $k_1^s(\omega_{31}) \approx 1\ \mu\text{m}^{-1}$. From Eqs. (4) and (11) for resonant optical transitions of rare-earth ions in crystals, e.g., Pr^{+3}-doped Y_2SiO_5 (demonstrated for EIT experiments [6]) with density $n \approx 10^{24}\ \text{m}^{-3}$, $\Gamma_{31} \cong 10^9$ rad/s, $\gamma_{21}^{-1} \gg 1\ \mu$s, $\varepsilon_1 \cong 1.3$, and $|\mathbf{d}| \approx -ea_0$ (with e the electron charge and a_0 the Bohr radius), we find $\alpha_0(\omega_{31}) \approx 10\ \mu\text{m}^{-1}$. Here we assume $z_0 \approx 1/k_1^s$ and $k_1^s \approx k_1^c$. Using Eq. (7) we compute

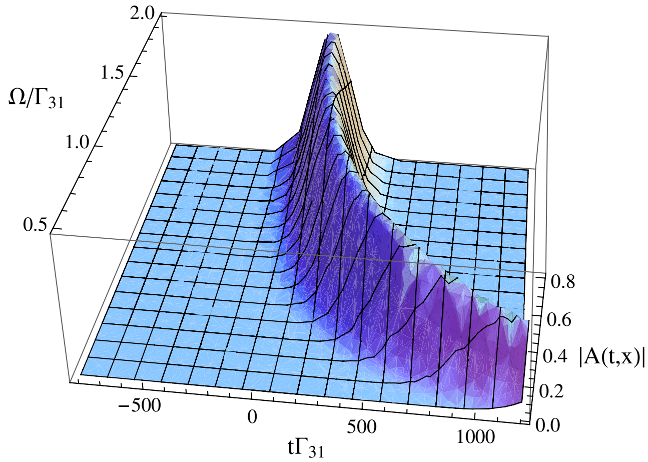

FIG. 4 (color online). 3D graph of pulse propagation profile as a function of time $t\Gamma_{31}$ and control field amplitude Ω/Γ_{31}, $\Gamma_{31} = 10^9$ rad/s, and $\kappa(\omega_{31}) = 0.01\kappa_0 = 100\ \mathrm{m}^{-1}$.

time delay and group velocity for a Gaussian amplitude envelope of the SP probe input pulse $\exp[-(t/\delta t)^2/2]$ in the medium at $x = 0$.

Input and output pulse profiles are presented in Figs. 4 and 5 for input pulse duration $\delta t = 100$ ns and for media lengths $x_1 = 1$ mm and $x_2 = 3$ mm ($< L = 1/\kappa$). As seen in Fig. 4, the time delay decreases with the control field amplitude as Ω^{-2}. Figure 5 shows the pulse profile for $\Omega = \Gamma_{31}$ when it propagates a distance $x_1 = 1$ mm and $x_2 = 3$ mm. The time delays are $2\delta t$ and $6\delta t$, respectively, whereas the amplitude has decreased only by factors 0.85 and 0.65, respectively. Thus the propagation length increased by more than $500/\alpha_0(\omega_{31})$ due to EIT of the SP field. Using these results we estimate the group velocity $v_g \approx 5000$ m/s and a compressed longitudinal envelope of the SP pulse $l_{SP} = v_g \delta t = 0.5$ mm, i.e., much smaller than the medium size. Thus the SP pulse can be successfully stored in the long-lived atomic coherence $\rho_{12}(t, x)$ and subsequently retrieved in accordance with an EIT quantum memory protocol [17].

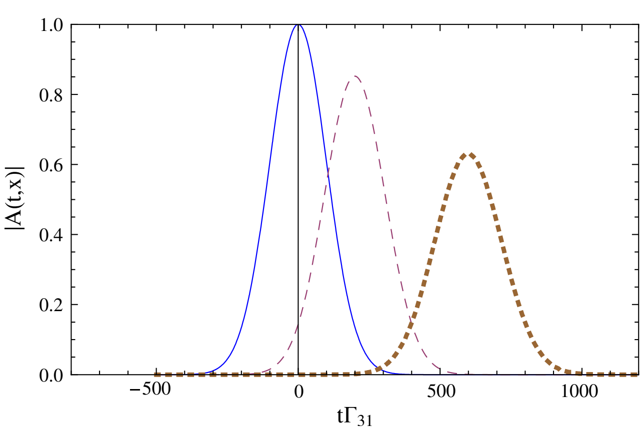

FIG. 5 (color online). The pulse propagation profile as a function of time $t\Gamma_{31}$ at different locations near the interface: blue (solid curve) at $x = 0$, red (dashed curve) at $x_1 = 1$ mm, and brown (dotted curve) at $x_2 = 3$ mm.

Conclusion.—We have demonstrated low loss SPs at the interface between two media with quite general electromagnetic properties, including dielectrics, metals, and metamaterials, and derived a closed form solution that provides deep insight into the system and control of spatially confined slow SP fields. We show that light pulses can be stored at the interface of two media exploiting EIT and low loss SP fields near a NIMM-dielectric interface.

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