

Near-Field Gain and Far-Field Control via a Plasmonic Time Crystal Slab

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Light-matter interactions can be substantially altered in the presence of time-varying media. We study the interaction between a harmonic electric dipole and a plasmonic time crystal slab. Temporal modulation of the plasma frequency enables near-field gain, allowing the dipole to absorb rather than emit energy, suppressing nonradiative losses. At the parametric resonance condition, the slab radiates strongly to the far field, producing 100% oscillations in the radiated power at distances up to 10^3 times the ϵ -near-zero wavelength. These findings reveal a new mechanism for controlling light-matter interaction in time-varying plasmonic systems.

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Introduction—Periodic temporal modulations of the optical properties of a material lift up the constraints of passivity and energy conservation, leading to the emergence of photonic time crystals (PTCs) [1]. These systems exhibit novel band structures with momentum gaps that host complex frequency modes [1–4] and have been shown to amplify both classical electromagnetic waves [5–9] and the quantum vacuum [10–14].

Additionally, PTCs can also fundamentally modify light-matter interactions, such as the emission properties of classical dipoles and quantum emitters, as shown in Refs. [6,15–20]. However, most of these works have neglected material dispersion and assumed a spatially infinite medium. Crucially, state-of-the-art platforms performing subcycle and order-unity temporal modulations of the electric permittivity, such as ϵ -near-zero (ENZ) materials with strong nonlinearities in the near-infrared [21–25], plasmonic metamaterials in the terahertz range [26], or transmission line metamaterials in the microwaves [8,27,28], are inevitably dispersive or confined to a finite region of space. Thus, including both of these complexities and understanding their roles is necessary to fully understand light-matter interactions in time-periodic media.

On the one hand, spatial confinement of PTCs to specific geometries, such as slabs [5,29–35] or spheres [36,37], shows that the lack of frequency conservation enables the different resonant modes of the structure to interact and hybridize, with the amplification phenomena displaying

features that depend on the spatial symmetries of the structure (planar, spherical, etc.). On the other hand, material dispersion allows one to engineer the size of the momentum gaps, which can be much larger in the dispersive case than in the nondispersive one [38,39]. Moreover, periodically modulated plasmonic media (that is, “plasmonic” time crystals) also support longitudinal modes. These can be amplified when the modulation frequency is resonant with the ENZ frequency, where $\text{Re}\{\epsilon(\omega_{\text{ENZ}})\} \sim 0$, giving rise to even larger momentum gaps than in the transverse case [40]. Therefore, accounting for spatial confinement and material dispersion in a PTC is not only of practical importance to model real-world scenarios, but it also allows for a higher tunability and enhancement of the amplification phenomena supported by these time-periodic systems.

In this Letter, we revisit the canonical setup for the study of nanoscale light-matter interactions, that is, dipole emission close to a nanostructure [41–43], but in the time-varying case. To this end, we study the emission properties of a harmonic electric dipole placed at a height h above a plasmonic time crystal slab of thickness d . We exploit the interplay between material dispersion, spatial structuring, and the time modulation to unveil new mechanisms for tailoring light-matter interactions both in the near field (h smaller than the ENZ wavelength, λ_{ENZ}) and the far field ($h \gg \lambda_{\text{ENZ}}$). We show that the time modulation introduces gain into the system, allowing the dipole to absorb energy in the near field. Therefore, the emitter ceases to be a source of electromagnetic radiation, instead becoming a sink. We also report very large (100% of the free-space value) and spatially long-lived oscillations of the radiated power in the far field. These oscillations arise from a sharp leaky resonance that is related to the Brewster condition and its negative frequency replica.

Physical model—Our system is a periodically time-modulated plasmonic slab, as seen in Fig. 1(a). The time

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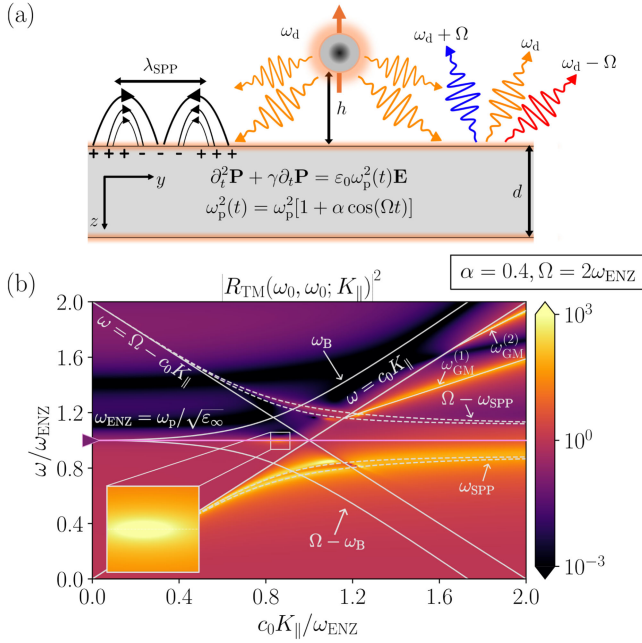


FIG. 1. (a) Diagram of the system under study: a harmonic electric dipole oscillating at frequency ω_d , located at a height h above a plasmonic time crystal slab. The slab supports surface plasmon polaritons. Additionally, owing to the lack of frequency conservation, waves can be reflected at different harmonics $\omega_d + n\Omega$, with n an integer. (b) Reflectance of the time-modulated metallic slab for transverse magnetic (TM) polarization vs frequency (normalized to the ENZ frequency $\omega_{\text{ENZ}} = \omega_p/\sqrt{\epsilon_\infty}$) and momentum (normalized to ω_{ENZ}/c_0). We consider values of $\alpha = 0.4$ and $\Omega = 2\omega_{\text{ENZ}}$. Inset: enlargement of the leaky resonance inside the light cone at $\omega \sim 0.5\Omega$.

modulation enters through the plasma frequency [9,44,45], which we assume is varied periodically in time as

$$\omega_p^2(t) = \omega_p^2[1 + \alpha \cos(\Omega t)]. \quad (1)$$

In the above, Ω is the modulation frequency and α is the modulation strength, while ω_p is the static (i.e., nonmodulated) plasma frequency. The conduction band electrons thus obey the following parametric-Drude model equation

$$\partial_t^2 \mathbf{P}(t) + \gamma \partial_t \mathbf{P}(t) = \epsilon_0 \omega_p^2(t) \mathbf{E}(t), \quad (2)$$

where $\mathbf{P}(t)$ is the polarization density of the conduction electrons, γ is their damping rate, $\mathbf{E}(t)$ is the externally applied electric field, and ϵ_0 the permittivity of free space.

The temporal periodicity of the plasma frequency allows us to use Floquet's theorem. Therefore, we make the following *Ansätze*: $P_j(t) = \sum_n P_{j,n} e^{-i\omega_n t}$ and $E_j(t) = \sum_n E_{j,n} e^{-i\omega_n t}$, where $j = x, y, z$ labels the different Cartesian components and $\omega_n = \omega + n\Omega$. Here, ω is a real frequency and $n \in \mathbb{Z}$. Under these Floquet *Ansätze*, we can explicitly solve the parametric-Drude differential

equation, with $P_{j,n} = \epsilon_0 \sum_{n'} \chi_{n,n'} E_{j,n'}$, where we introduce the following Floquet susceptibility matrix:

$$\chi_{n,n'} = \left(\epsilon_\infty - 1 - \frac{\omega_p^2}{\omega_n^2 + i\gamma\omega_n} \right) \delta_{n,n'} - \frac{\alpha}{2} \frac{\omega_p^2}{\omega_n^2 + i\gamma\omega_n} (\delta_{n+1,n'} + \delta_{n-1,n'}). \quad (3)$$

The first line in Eq. (3) corresponds to the static response of the system, while the second line contains the effects of the temporal modulation. Additionally, the background permittivity ϵ_∞ (not captured by the driven-Drude equation, but still demanded by the Kramers-Kronig relations [46]) has been included. The material response is thus dispersive, i.e., frequency dependent, and time modulated, since the frequency domain susceptibility is not a scalar, but a matrix. Hence, the time-modulated linear response allows different harmonics to interact.

Next, we solve Maxwell equations $\nabla \times \mathbf{E} = -\partial_t \mathbf{B}$ and $\nabla \times \mathbf{H} = \partial_t \mathbf{D}$, making use of the constitutive relations $B_{j,n} = \mu_0 H_{j,n}$ and $D_{j,n} = \epsilon_0 \sum_{n'} \epsilon_{n,n'} E_{j,n'}$, where $\epsilon_{n,n'} = \delta_{n,n'} + \chi_{n,n'}$ is the Floquet permittivity matrix. By imposing the boundary conditions at the slab interface for all harmonics (see Supplemental Material [47]), we calculate the reflected field for given polarization (with σ standing for transverse electric and transverse magnetic waves) and lateral momentum $K_{||}$, both of which are conserved quantities. Frequency, on the other hand, is only conserved modulo Ω and, therefore, the reflected field will contain several harmonics $\omega_n = \omega + n\Omega$ [34,36]. Therefore, the tangential component of the reflected field at the air-metal interface (for given σ and $K_{||}$) reads $E_\sigma^{(\text{ref})}(t; K_{||}) = \sum_n E_\sigma^{(\text{ref})}(\omega_n, K_{||}) e^{-i\omega_n t}$, where

$$E_\sigma^{(\text{ref})}(\omega_n, K_{||}) = \sum_{n'} R_\sigma(\omega_n, \omega_{n'}; K_{||}) E_\sigma^{(\text{in})}(\omega_{n'}, K_{||}), \quad (4)$$

with $E_\sigma^{(\text{in})}(\omega_{n'}, K_{||})$ being the amplitudes of the tangential incident field.

We have introduced generalized Fresnel-Floquet reflection coefficients $R_\sigma(\omega_n, \omega_{n'}; K_{||})$ that couple an incident frequency $\omega_{n'}$ with an outgoing one ω_n , where $\omega_n - \omega_{n'} = (n - n')\Omega$. Since we shall deal with a harmonic electric dipole, which oscillates at a single frequency, we only discuss in here the frequency-conserving coefficient $R_\sigma(\omega_0, \omega_0; K_{||})$ and for $\sigma = \text{TM}$ in order to maximize dipole coupling. Despite conserving the frequency of the incident wave, $R_{\text{TM}}(\omega_0, \omega_0; K_{||})$ contains the effects of the different frequency bands $\omega_n = \omega + n\Omega$, which must all be taken into account for an accurate computation of $R_{\text{TM}}(\omega_0, \omega_0; K_{||})$ [47].

Fresnel-Floquet reflection coefficient—In order to facilitate a quantitative discussion, we assume the slab to be

made out of indium tin oxide (ITO), an ENZ material with a strong nonlinear response [22]. The optical response of ITO is dominated by the conduction band electrons and, therefore, is properly described with a Drude model [22,48]. The values for the Drude-model parameters and the slab's thickness are all taken from experimental data, in particular, those found in Ref. [44], with plasma wavelength $\lambda_p = 2\pi c_0/\omega_p = 654$ nm, damping rate $\gamma = 0.13$ (fs) $^{-1}$, background permittivity $\epsilon_\infty = 4$, and thickness $d = 310$ nm, although our results are not specific to this set of parameters. In Fig. 1(b), we plot the reflectance of the time-modulated metallic slab for TM-polarized waves, $|R_{\text{TM}}(\omega_0, \omega_0; K_\parallel)|^2$, with $\alpha = 0.4$ and $\Omega = 2\omega_{\text{ENZ}}$. Additionally, we highlight both the light cone (solid gray line) and the ENZ frequency (horizontal purple line), which separates the metallic ($\omega < \omega_{\text{ENZ}}$) and insulating regimes ($\omega > \omega_{\text{ENZ}}$). We also mark with dashed gray lines the surface plasmon polaritons (SPPs) supported by the metallic slab ω_{SPP} [49–54]; these modes are confined to the air-metal interface, propagate along it [52,55], and appear as poles of the reflection coefficient for TM-polarized waves, thus leading to diverging values of the reflectivity [54,56].

For $K_\parallel$ large enough, we reach the quasistatic regime, in which the two SPPs converge to the surface plasmon (SP) [49], whose frequency is $\omega_{\text{SP}} = \omega_p/\sqrt{1 + \epsilon_\infty}$. The latter gives a flat ($K_\parallel$ -independent) resonance which, in the static case, greatly enhances the quenching of power radiated by a dipole placed in the near field of the slab [57,58]. Additionally, above the ENZ frequency we can see the first two guided modes (GMs) of the slab [46], $\omega_{\text{GM}}^{(1)}$ and $\omega_{\text{GM}}^{(2)}$ highlighted in solid gray, as well as the Brewster condition ω_B , also marked in solid gray, at which the reflection coefficient for TM polarization vanishes.

We also highlight the upshifted negative frequency replicas of the SPPs, $\Omega - \omega_{\text{SPP}}$, which are seen to interact with ω_B , and mark as well the upshifted negative frequency replica of the Brewster condition, $\Omega - \omega_B$. Surprisingly, the temporal modulation makes both the reflectivity poles and zeros of the static slab to interact, with the zeros effectively behaving as modes of the system. This is indeed the case for ω_B and $\Omega - \omega_B$, which interact at $\omega = 0.5\Omega = \omega_{\text{ENZ}}$ and $K_\parallel = 0$. This interaction results in the formation of a sharp leaky resonance inside the light cone at $\omega = \omega_{\text{ENZ}}$, which appears as a peak in the inset of Fig. 1(b). The resonance is first formed at $K_\parallel = 0$ and moves toward larger values of $K_\parallel$ as α increases, resembling the behavior of the edge of a $K_\parallel$ gap (a detailed account of this can be found in Supplemental Material [47]).

Dipole radiation—We now place a harmonic electric dipole at a height h above the plasmonic slab, as shown in Fig. 1(a). The dipole oscillates at a frequency ω_d , and its dipole moment $\mathbf{d}$ is assumed to be perpendicular to the air-metal interface to maximize the coupling to TM modes. The energy exchanged per unit time between the dipole and

the electric field—which includes both the dipole's own field and the one reflected from the surface—is given by $P(t) = \int d^3x \mathbf{J}_d \cdot \mathbf{E}$, where $\mathbf{J}_d = -i\omega_d \mathbf{d} e^{-i\omega_d t} \delta(\mathbf{x} - \mathbf{x}_0) + \text{c.c.}$ is the current associated with the emitter. The field reflected from the plasmonic slab contains many different harmonics $\omega_d + n\Omega$, as seen in Fig. 1. However, as the dipole is monochromatic, for $t \gg 1/\omega_d$, only the fundamental one ($n = 0$) has a nonvanishing contribution to the total radiated power $P_{\text{tot}} = \lim_{t \rightarrow \infty} [1/(t - t_0)] \int_{t_0}^t dt' P(t')$ [47,54], where $t_0 \leq t$ is arbitrary and determines the initial time at which the dipole starts radiating. The total power can be separated into its radiative and nonradiative contributions, with $P_{\text{tot}} = P_{\text{rad}} + P_{\text{nonrad}}$. Particularizing the latter two quantities for $\mathbf{d}$ perpendicular to the surface we have

$$\frac{P_{\text{rad}}}{P_0} = 1 + \frac{3c_0^3}{2\omega_d^3} \int_0^{\omega_d/c_0} \frac{dK_\parallel K_\parallel^3}{\kappa_z} \text{Re}\{R_{\text{TM}}(\omega_d, \omega_d; K_\parallel) e^{-i2k_z h}\} \quad (5)$$

and

$$\frac{P_{\text{nonrad}}}{P_0} = \frac{3c_0^3}{2\omega_d^3} \int_{\omega_d/c_0}^\infty \frac{dK_\parallel K_\parallel^3}{\kappa_z} \text{Im}\{R_{\text{TM}}(\omega_d, \omega_d; K_\parallel)\} e^{-2\kappa_z h}. \quad (6)$$

In Eqs. (5) and (6), $P_0 = \mu_0 |\mathbf{d}|^2 \omega_d^4 / (3\pi c_0^3)$ is the total radiated power in free space, which is given by Larmor's formula and to which we normalize P_{rad} and P_{nonrad} . On the other hand, $k_z = \sqrt{(\omega_d/c_0)^2 - K_\parallel^2}$ and $\kappa_z = \sqrt{K_\parallel^2 - (\omega_d/c_0)^2}$ are the wave numbers for plane and evanescent waves along the z axis, which we take to be perpendicular to the air-metal interface (see Fig. 1). Thus, P_{rad} is seen to be completely determined by the modes inside the light cone, while P_{nonrad} is given by the ones outside of it, such as the GMs or the SPPs.

Far-field control—In Fig. 2, we plot the power P vs the dipole's height above the surface h (normalized to the ENZ wavelength $\lambda_{\text{ENZ}} = \sqrt{\epsilon_\infty} \lambda_p$) and for $\omega_d \sim 0.5\Omega$. We consider modulation parameters of $\alpha = 0.4$ and $\Omega = 2\omega_{\text{ENZ}}$, as in Fig. 1(b). The solid black line in Fig. 2 corresponds to the total power emitted in the time-modulated case P_{tot} , while the solid gray line is the total power emitted in the static case $P_{\text{tot}, \alpha=0}$. We also plot P_{rad} in dashed green and P_{nonrad} in solid blue, which is enhanced with respect to the static case at small emitter-slab distances. The power at large emitter-slab distances shows unexpected features, with P_{tot} displaying large and spatially long-lived oscillations that are a 100% of the free-space value and which are completely determined by the radiative contribution to the emitted power. These far-field oscillations stem from

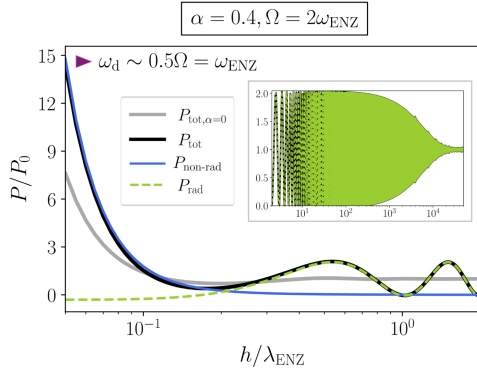


FIG. 2. Power emitted by a dipole oscillating at $\omega_d \sim 0.5\Omega = \omega_{\text{ENZ}}$ as a function of its height above the surface. Inset (P/P_0 vs h/λ_{ENZ}): the spatial extent of the far-field oscillations. Modulation parameters are set to $\alpha = 0.4$ and $\Omega = 2\omega_{\text{ENZ}}$, as in Fig. 1.

the sharp peak inside the light cone at $\omega = \omega_{\text{ENZ}}$ [see inset of Fig. 1(b)], which is rooted in the interaction between ω_B and $\Omega - \omega_B$.

In this way, the emission of light by the dipole in the far field can be enhanced by a factor of 2 or be completely suppressed, depending on the value of h , and at distances of up to $h \sim 10^3 \lambda_{\text{ENZ}}$. Thus, the time-modulated plasmonic slab acts as an antenna at $\omega \sim 0.5\Omega$, radiating electromagnetic energy to the far-field and modifying the emission of light at distances much larger than in the static case. These oscillations become compressed as h increases and eventually die out for $h \sim 5 \times 10^4 \lambda_{\text{ENZ}}$ (see the inset of Fig. 2), where the free-space value is recovered.

Critically, this leaky resonance arising when the plasmonic time crystal slab is driven parametrically is the signature of an instability. In this unstable regime, the system's scattering matrix has poles above the real frequency axis (that is, $\text{Im}\{\omega_{\text{pole}}\} > 0$) [56], so that the slab supports resonant modes that are exponentially amplified in time and exhibits lasing [20,28,47,59]. Importantly, in our calculation of the total power we only account for the field generated by the dipole and the field reflected from the surface, which are dominant for short times $t \ll 1/\max(\text{Im}\{\omega_{\text{pole}}\})$. We thus neglect the fields of the exponentially growing resonant modes, whose contribution to the total power obscures those of Eqs. (5) and (6) at long times, $t \sim 1/\max(\text{Im}\{\omega_{\text{pole}}\})$. Crucially, nonlinearities of the material response eventually lead to the exponentially amplified modes to saturate [28,60].

Near-field gain—A glance at the formula for P_{nonrad} , Eq. (6), shows that the integrand's sign is completely determined by the imaginary part of the generalized reflection coefficient at the dipole's frequency $\text{Im}\{R_{\text{TM}}(\omega_d, \omega_d; K_{\parallel})\}$. Therefore, a negative value of $\text{Im}\{R_{\text{TM}}(\omega_d, \omega_d; K_{\parallel})\}$ yields a negative contribution P_{nonrad} . This can be understood as the dipole absorbing energy instead of radiating it, with the electromagnetic field doing

work on the dipole instead of the opposite. Bearing this in mind, we can further separate P_{nonrad} into a loss contribution P_{loss} , determined by the $\text{Im}\{R_{\text{TM}}(\omega_d, \omega_d; K_{\parallel})\} \geq 0$ regions, and a gain one P_{gain} , given by the $\text{Im}\{R_{\text{TM}}(\omega_d, \omega_d; K_{\parallel})\} < 0$ regions. Then, the nonradiative power is $P_{\text{nonrad}} = P_{\text{loss}} - P_{\text{gain}}$, where P_{gain} is defined to be positive [19,61]. Finally, summing radiative, loss, and gain contributions, the total power reads $P_{\text{tot}} = P_{\text{rad}} + P_{\text{loss}} - P_{\text{gain}}$.

Notice that, in passive media, causality imposes $P_{\text{nonrad}} \geq 0$ (and $P_{\text{gain}} = 0$) as $\text{Im}\{R_{\sigma}^{(0)}(\omega; K_{\parallel})\} \geq 0$ outside the light cone, where $R_{\sigma}^{(0)}(\omega; K_{\parallel})$ is the reflection coefficient in the static case. In media with gain, this restriction is lifted, and it thus becomes possible to have $\text{Im}\{R_{\sigma}(\omega_0, \omega_0; K_{\parallel})\} < 0$ for $K_{\parallel} \geq \omega/c_0$ and, consequently, $P_{\text{gain}} > 0$.

In Fig. 3(a), we plot $\text{Im}\{R_{\text{TM}}(\omega_0, \omega_0; K_{\parallel})\}$ as a function of the frequency ω and lateral momentum $K_{\parallel}$, in this case for $\alpha = 0.25$ and $\Omega = 1.90\omega_{\text{ENZ}}$; blue-colored regions have $\text{Im}\{R_{\text{TM}}(\omega_0, \omega_0; K_{\parallel})\} > 0$, while red-colored ones have

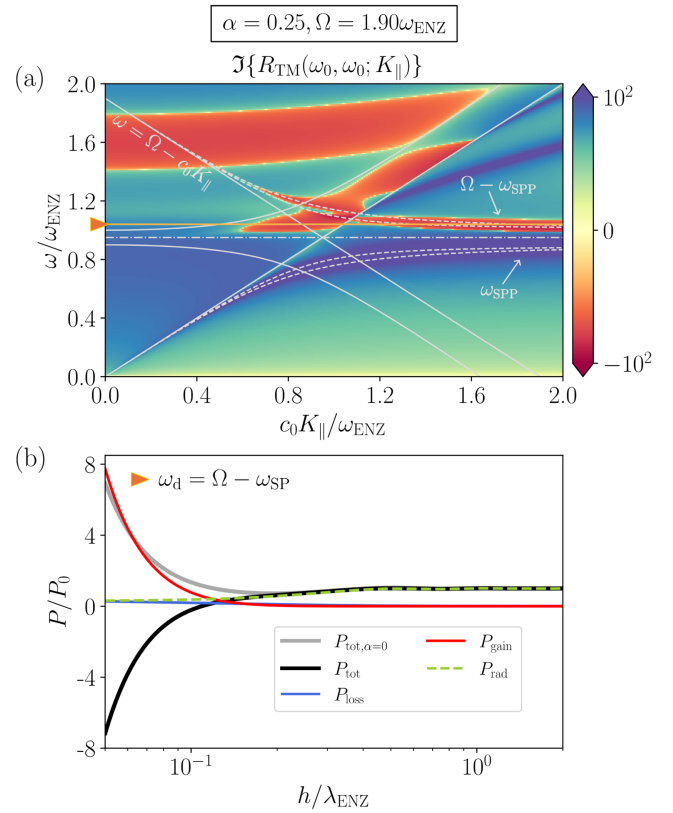


FIG. 3. (a) Imaginary part of the frequency-conserving TM reflection coefficient vs frequency and momentum. Outside the light cone, blue-colored regions correspond to near-field losses, while red-colored ones are associated to near-field gain. (b) Total power emitted by a dipole oscillating at $\omega_d = \Omega - \omega_{\text{SP}}$ as a function of its height above the surface. Gain dominates the near field, in which the emitted power becomes negative. Modulation parameters are set to $\alpha = 0.25$ and $\Omega = 1.90\omega_{\text{ENZ}}$.

$\text{Im}\{R_{\text{TM}}(\omega_0, \omega_0; K_{\parallel})\} < 0$. According to the preceding discussion, we can identify blue regions outside the light cone with loss and red ones with gain.

A glance at Fig. 3(a) shows that $\text{Im}\{R_{\text{TM}}(\omega_0, \omega_0; K_{\parallel})\}$ takes negative values along the $\Omega - \omega_{\text{SPP}}$'s dispersion curves, which appear as red colored in the figure. The SP replica (the quasistatic limit of $\Omega - \omega_{\text{SPP}}$) is particularly interesting, since a dipole whose frequency is matched to $\Omega - \omega_{\text{SP}}$ will couple to a flat region of gain, absorbing energy instead of radiating it. More importantly, the closer the emitter is to the surface, the more it will absorb. However, such an increase in dipole absorption will eventually saturate when the intrinsic anharmonicities of the emitter become dominant.

In Fig. 3(b), we plot the power emitted by a dipole oscillating at $\omega_d = \Omega - \omega_{\text{SP}}$, in this case, $\alpha = 0.25$ and $\Omega = 1.90\omega_{\text{ENZ}}$. Crucially, for these modulation parameters, the system is in a stable regime, unlike in the results of Figs. 1(b) and 2. We see that P_{loss} (solid blue curve) vanishes, while P_{gain} (solid red curve) indeed becomes increasingly large as h becomes smaller. Consequently, P_{tot} decreases as the air-metal interface is approached, eventually becoming negative in the near field and reaching values of $P_{\text{tot}} \sim -7P_0$ for $h = 0.05\lambda_{\text{ENZ}}$.

Crucially, it is at the transition from the SPP-replica dominating regime to the quasistatic one that P_{gain} begins to grow large and, correspondingly, that P_{tot} becomes negative, as shown in Supplemental Material [47]. Therefore, the quenching of the radiated power at ω_{SP} is transformed by the time modulation into nonradiative gain at $\Omega - \omega_{\text{SP}}$, with nonradiative losses becoming completely suppressed in the process. Additionally, as opposed to the leaky resonance allowing the control of the far field, the near-field gain is achieved without relying on an instability. Indeed, the gain solely originates from the negative frequency replica of the SP [47].

Finally, we comment on the specific values chosen for the modulation parameters: for the case of far-field control, the system must be in the unstable regime, for which α and Ω ought to be large enough that the poles of the scattering matrix cross the real frequency axis. In particular, having $\Omega \sim 2\omega_{\text{ENZ}}$ makes ω_B and $\Omega - \omega_B$ interact more strongly for a given α , since in that case the interaction point lies at $\omega \sim \omega_{\text{ENZ}}$ and $K_{\parallel} \sim 0$, where both modes are flat [47]. On the other hand, for the system to support near-field gain, it is sufficient to have $\Omega \gtrsim 2\omega_{\text{SP}}$, so that $\Omega - \omega_{\text{SP}}$ lies close to the light cone. Critically, α does not need to be as large as for the previous case, since the system can be kept within its stable regime. However, it does need to be above some threshold value determined by the losses γ in order to obtain gain from the SPP replica.

Conclusions—We have demonstrated that temporal modulation of a plasmonic slab can fundamentally alter dipole emission, transforming near-field quenching into nonradiative gain and enabling far-field radiation control at

unprecedented distances. In particular, coupling to the negative frequency replica of the surface plasmon produces a flat gain band that turns a dipole into an energy absorber, while the emergence of a leaky resonance associated with the Brewster condition and its negative frequency replica yields large, long-lived far-field oscillations comparable to the free-space emission level and up to distances of $10^3\lambda_{\text{ENZ}}$. These results establish a new route for tailoring light-matter interaction in time-varying plasmonic systems by simultaneously exploiting spatial confinement and material dispersion. Beyond the classical regime studied here, our findings suggest analogous phenomena for quantum emitters [6,17], opening opportunities for sustained excitation and emission engineering in plasmonic time crystals.

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Data availability—The data that support the findings of this article are openly available [62].

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