Reflective optical limiter based on resonant transmission

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Optical limiters transmit low-level radiation while blocking electromagnetic pulses with excessively high energy (energy limiters) or with excessively high peak intensity (power limiters). A typical optical limiter absorbs most of the high-level radiation, which can cause its overheating and destruction. Here we introduce the concept of a reflective energy limiter which blocks electromagnetic pulses with excessively high total energy by reflecting them back to space, rather than absorbing them. The idea is to use a defect layer with temperature-dependent loss tangent embedded in a low-loss photonic structure. The low-energy pulses with central frequency close to that of the localized defect mode will pass through. But if the cumulative energy carried by the pulse exceeds certain level, the entire photonic structure becomes highly reflective (not absorptive) within a broad frequency range. The underlying physical mechanism is based on self-regulated impedance mismatch which increases dramatically with the cumulative energy carried by the pulse.

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I. INTRODUCTION

The protection of photosensitive optical components from high incident radiation has applications ranging from microwave and optical communications to optical sensing [1–3]. As a result, a considerable research effort has focused on developing protection schemes and materials that provide control of high-level optical and microwave radiation and prevent damages of optical sensors (including the human eye) and microwave antennas [4–9]. Optical limiters constitute an important class of such protection devices. They are supposed to transmit low-level radiation, while blocking light pulses with excessively high level of radiation. A typical passive optical limiter absorbs most of the high-level radiation, which can cause its overheating and destruction. The most common realization of a passive optical limiter is provided by a single protective layer with complex permittivity $\epsilon = \epsilon' + i\epsilon''$, where the imaginary part ϵ'' increases sharply with the radiation level. For low-level radiation, the absorption is negligible and the protective layer is transparent. An increase in the radiation level results in an increase in ϵ'', which renders the protective layer opaque. As a consequence, most of the high-level radiation will be absorbed by the limiter. If the same protective layer is incorporated into a certain photonic layered structure, the entire multilayer can become highly reflective for high-level radiation, while remaining transmissive at certain frequencies if the radiation level is low. Such a photonic reflective limiter can be immune to overheating and destruction by high-level laser radiation, which is our main objective.

The physical reasons for the sharp increase in ϵ'' with the radiation level can be different. For instance, it can be photoconductivity, two-photon absorption, heating, or any combination of the above mechanisms. In our previous publication [10] we considered the particular case of a strong nonlinear dependence of ϵ'' of the protective layer on light intensity. This can be attributed, for instance, to a two-photon absorption. We showed that incorporation of such a nonlinear layer in a properly designed low-loss layered structure makes the entire assembly act as a reflective power limiter. In this paper, we consider a more practical particular case where the increase in ϵ'' is due to heating of the protective layer. We show that, depending on the pulse duration as compared to the thermal relaxation time, the properly design layered structure incorporating such a protective layer can act as a reflective energy limiter or as a reflective power limiter. Specifically, for short pulses, such a layered structure acts as an energy limiter, reflecting light pulses carrying excessively high energy. By comparison, for sufficiently long pulses, the same layered structure will act as a power limiter. In either case, most of the incident radiation will be reflected back to space, even though a standalone protective layer would act as an absorptive optical limiter.

The proposed architecture consists of a (protective) defect layer embedded in a low-loss Bragg grating. In contrast to the reflective power limiter introduced in Ref. [10], the defect layer does not have to be nonlinear, but it must display strong temperature dependence $\epsilon''(T)$ of the imaginary part of its permittivity. If the total energy carried by the pulse is low, $\epsilon''(T)$ remains small enough to support a localized mode and the resonant transmittance associated with this mode. If, on the other hand, the energy carried by the pulse exceeds a certain level, the defect layer becomes lossy enough to suppress the localized mode, along with the resonant transmittance. The entire stack turns highly reflective, which is consistent with our goal. We refer to this limiter as a *reflective energy limiter* in order to distinguish it from the nonlinear reflective power limiter introduced in Ref. [10]. Finally, if the pulse duration significantly exceeds the thermal relaxation time of the defect layer, the entire layered structure will again act as a reflective power limiter with the cutoff light intensity determined by the thermal relaxation time of the defect layer, not by the nonlinearity in ϵ'', as was the case in Ref. [10].

The organization of the paper is as follows. In Sec. II we clarify the different mechanisms underlying a reflective energy limiter (the theme of the present study) and a reflective power limiter (the theme of Ref. [10]). In Sec. III, a conceptual design

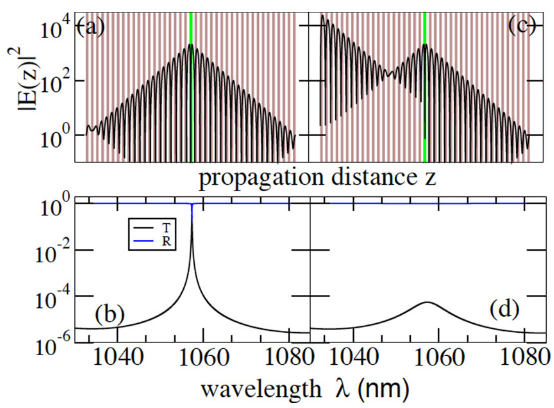

FIG. 1. (Color online) A schematics of a reflective energy limiter. Two identical lossless Bragg reflectors are placed on the left and right of a lossy layer (green [gray]). The value of ϵ'' in the defect layer is an increasing function of temperature. (a) Field distribution at the frequency ω_r of resonant transmission for an incident pulse with low energy—the field amplitude at the location of the defect layer is exponentially higher than that of the incident wave. (b) Transmittance vs light wavelength for low incident light energy. (c) Field distribution at the frequency of maximum transmittance for an incident pulse with high energy—the amplitude of the suppressed localized mode is lower than that of the incident wave. (d) Transmittance vs wavelength for an incident pulse with high energy.

for the reflective energy limiter is presented, along with the mathematical formalism used in our calculations. In Sec. IV, we analyze the role of thermal conductivity. The latter plays an important role if the pulse duration is comparable or exceeds the thermal relaxation time of the defect layer. Our conclusions are given in Sec. V.

II. BASIC CONCEPTS: REFLECTIVE ENERGY LIMITER VERSUS REFLECTIVE POWER LIMITER

Before proceeding with our analysis we would like to clarify the two notions of reflective energy limiter and reflective power limiter. In both cases, the limiter structure consists of a defect layer with complex permitivity $\epsilon_d = \epsilon'_d + i\epsilon''_d$ which is embedded in a Bragg grating; see Fig. 1. However, as we will explain, the basic principle behind the limiting action is completely different for these two type of limiters.

Let us start with the power limiter considered in our previous publication [10]. In this case, the key assumption was that, at any moment t, the complex permittivity $\epsilon(t)$ of the defect layer is a function of the instantaneous value of the oscillating electric field $E(t)$. It also implies that the permittivity $\epsilon(t)$ is assumed independent of the field intensity $E(t')$ at $t' < t$. This is a standard assumption in nonlinear optics, but it fails if heating or some other time-cumulative effects are significant. Still, the very definition of power limiter adopted in Ref. [10] is based on this assumption. In Ref. [10] we have considered, as an example case, that the main mechanism behind $\epsilon''_d(E)$ is a two-photon absorption process i.e., $\epsilon''_d(E) = \chi|E|^2$.

Let us turn from the power limiter to the energy limiter considered in this paper. The key assumption now is that, at any moment t, the field intensity $E(t)$ is too small to affect the instantaneous value $\epsilon(t)$ of the complex permittivity of the defect layer. In this sense, our system is linear. At the same time, if the pulse duration is long enough, it can cause heating, or some other time-cumulative effects resulting in a gradual change of ϵ in time. The relative change in permittivity during the oscillation period should be negligible, which is a very realistic assumption. This approximation is often referred to as adiabatic approximation. If all the heat released in the defect layer stays there, then the system can act as an energy limiter, because its transmittance is a function of the total energy carried by the light pulse, rather than just its intensity. Finite thermal conductivity can affect the result, and we address this issue too in our analysis below. But in any event, the transmission characteristic of such a system will depend on pulse intensity and duration. In this respect, it does not behave as a simple linear system.

To summarize, in Ref. [10] we assumed a strong instantaneous nonlinearity of the defect layer, but no time-cumulative effects (like heating). This is why the optical limiter in Ref. [10] is only sensitive to the pulse peak intensity and not to its duration. In contrast, in this paper we assume a negligible instantaneous nonlinearity, while the heating is essential for the performance of the limiter. In practice, there might be a combination of the two mechanisms. The bottom line, though, is that no matter what causes the rise in ϵ'', the layered structure in Fig. 1 will act as a reflective optical limiter. In contrast, a standalone defect layer would absorb most of the high-level radiation.

We will demonstrate the concept of energy limiter, highlighted above, using a simplified model. In our modeling we will omit any dispersive phenomena of ϵ_d originating from the material considered (temporal dispersion). We did this for a reason: Indeed, the change in ϵ_d due to heating required for the limiter to perform is usually at least two or three orders of magnitude, which is much greater compared to the typical temporal dispersion of optical materials. We will also assume that the pulse duration is much larger than the carrier period. This justifies the use of the adiabatic approximation, which means that the heat release during one carrier period of oscillation is infinitesimally small. Finally we have assumed a simplified dependence of ϵ''_d from the temperature of the defect layer. More realistic schemes or dependences will only mask the demonstration of the concept with unnecessary numerical complications.

III. PHYSICAL STRUCTURE AND MATHEMATICAL MODEL

We consider two identical lossless Bragg reflectors consisting of two alternating layers. Each mirror consists of forty layers which are placed at $-L \leqslant z \leqslant 0$ and $d \leqslant z \leqslant L + d$. For the sake of the discussion we assume that the layers consist of Al_2O_3 and SiO_2 with corresponding permittivities $\epsilon_1 = 3.08$ and $\epsilon_2 = 2.1$. These values are typical for these materials at wavelengths $\lambda \sim 1$ μm. The width of layers is assumed to be $d_1 = 151$ nm and $d_2 \approx 183$ nm respectively. At $0 \leqslant z \leqslant d$ we introduce a defect lossy layer with complex permittivity $\epsilon_\mathrm{d} = \epsilon'_\mathrm{d} + i\epsilon''_\mathrm{d}$. We further assume that the imaginary part of the permittivity of the defect layer depends on the temperature T, i.e., $\epsilon''_\mathrm{d} = \epsilon''_\mathrm{d}(T)$. For simplicity, we assume linear dependence, i.e., $\epsilon''_\mathrm{d}(T) = c_1 + c_2 T$, where c_1,c_2 are some characteristic constants of the defect. Below we assume that $\epsilon'_{\rm d} = 12.11$ (which is a typical value for, say, GaAs at near infrared), $c_1 = 10^{-5}$ and $c_2 = 1$ while the width of the defect layer is taken to be $d = 151$ nm.

The transport characteristics $\mathcal{T},\mathcal{R},\mathcal{A}$ of our setup and the field profile at any frequency can be calculated via the transfer matrix approach. Specifically, a time-harmonic electric field of frequency ω satisfies the Helmholtz equation:

$$\frac{\partial^2 E(z)}{\partial z^2} + \frac{\omega^2}{c^2}\epsilon(z)E(z) = 0. \tag{1}$$

At each layer inside the grating, Eq. (1) admits the solution $E^{(j)} = E_f^{(j)} \exp(in_jkz) + E_b^{(j)} \exp(-in_jkz)$, where $n_j = \sqrt{\epsilon_j}$ is the refraction index of the jth layer and k is the wave vector $k = \omega/n_0c$ (c is the speed of light in the vacuum and n_0 is the refractive index of air). Imposing continuity of the field and its derivative at each layer interface, as well as taking into consideration the free propagation in each layer, we get the following iteration relation:

$$\begin{pmatrix} E_f^{(j)} \\ E_b^{(j)} \end{pmatrix} = \mathcal{M}^{(j)} \begin{pmatrix} E_f^{(j-1)} \\ E_b^{(j-1)} \end{pmatrix}; \quad \mathcal{M}^{(j)} = P_R^{(j)} Q^{(j)} K^{(j)} P_L^{(j)}, \tag{2}$$

where

$$Q^{(j)} = \begin{pmatrix} e^{ikn_jd_j} & 0 \\ 0 & e^{-ikn_jd_j} \end{pmatrix},$$
$$K^{(j)} = \begin{pmatrix} \frac{n_j+n_{j-1}}{2n_j} & \frac{n_j-n_{j-1}}{2n_j} \\ \frac{n_j-n_{j-1}}{2n_j} & \frac{n_j+n_{j-1}}{2n_j} \end{pmatrix},$$
$$P_R^{(j)} = \begin{pmatrix} e^{-ikn_jz} & 0 \\ 0 & e^{ikn_jz} \end{pmatrix},$$
$$P_L^{(j)} = \begin{pmatrix} e^{ikn_{j-1}(z-d_j)} & 0 \\ 0 & e^{-ikn_{j-1}(z-d_j)} \end{pmatrix}. \tag{3}$$

At the same time the field outside the layered structured can be written as $E_0^-(z) = E_f^- \exp(ikz) + E_b^- \exp(-ikz)$ for $z < -L$ and $E_0^+(z) = E_f^+ \exp(ikz) + E_b^+ \exp(-ikz)$ for $z > L + d$. The amplitudes of forward- and backward-propagating waves on the left $z < -L$ and right $z > L + d$ domains are related via the total transfer matrix $\mathcal{M} = P_R^{(2N+2)} K^{(2N+2)} \Pi_j \mathcal{M}^{(j)}$ (where N is the number of layers on each grating and $n_{2N+2} = n_0$):

$$\begin{pmatrix} E_f^+ \\ E_b^+ \end{pmatrix} = \begin{pmatrix} \mathcal{M}_{11} & \mathcal{M}_{12} \\ \mathcal{M}_{21} & \mathcal{M}_{22} \end{pmatrix} \begin{pmatrix} E_f^- \\ E_b^- \end{pmatrix}. \tag{4}$$

The transmittance and reflectance and the field profile, say for a left incident wave, can be obtained by iterating backwards Eqs. (2) and (4) together with the boundary conditions $E_b^+ = 0$ and $|E_f^+| = 1$ (due to the linearity of the equations, one can always impose a value for the outgoing field and calculate via a backward iteration of the transfer matrices the corresponding input field [11]). Specifically we have $\mathcal{T} \equiv |E_f^+/E_f^-|^2$; $\mathcal{R} \equiv |E_b^-/E_f^-|^2$. These can be expressed in terms of the transfer matrix elements as $\mathcal{T} = |\frac{1}{\mathcal{M}_{22}}|^2$; $\mathcal{R} = |\frac{\mathcal{M}_{21}}{\mathcal{M}_{22}}|^2$. The absorption coefficient $\mathcal{A}$ can then be evaluated in terms of transmittances and reflectances as $\mathcal{A} \equiv 1 - \mathcal{T} - \mathcal{R}$.

IV. THEORETICAL ANALYSIS

In the case that the permittivity of the defect layer is replaced by $\epsilon_d = \epsilon_1$, the whole structure is periodic and displays a typical dispersion relation consisting of transparent frequency windows (bands) where light is transmitted with near-unity transmittance alternated with frequency windows (gaps) where the incident light is experiencing almost complete reflection.

When the defect is included in the middle of the grating, for zero temperature $T = 0$ corresponding to permittivity $\epsilon_{\rm d} \approx \epsilon'_{\rm d}$, the layered structure supports a localized resonant defect mode [see Fig. 1(a)] with a frequency lying in a photonic band gap of the Bragg grating [see Fig. 1(b)]. For the specific setup that we consider here, we find that a resonant mode is located in the middle of the gap at wavelength $\lambda_r \approx 1060$ nm. This defect mode is localized in the vicinity of the defect layer and decays exponentially away from the defect [see Fig. 1(a)]. In the vicinity of the localized mode frequency ω_r, the entire layered structure displays a strong resonant transmission due to the excitation of the localized mode [see Fig. 1(b)]. In other words, the transmittance is $\mathcal{T}(\omega_r) \equiv \mathcal{T}_r \approx 1$ while the reflectance and the absorption in the absence of any losses are $\mathcal{R}(\omega_r) \equiv \mathcal{R}_r \approx 0$ and $\mathcal{A}(\omega_r) \equiv \mathcal{A}_r \approx 0$ respectively. This picture is still applicable even in the presence of small (but nonzero) dissipative permittivity $\epsilon''_{\rm d} \neq 0$ [see Figs. 1(a) and 1(b)].

An alternative expression for the absorption coefficient $\mathcal{A}$ can be given in terms of the permittivity and the field intensity $|E(z)|^2$ inside the defect layer. The resulting expression is derived by subtracting the product of Eq. (1) with $E^*(z)$ from its complex conjugate form and then integrating the outcome over the interval $-L \leqslant z \leqslant L$. We get

$$\left(E^*\frac{dE}{dz} - E\frac{dE^*}{dz}\right)_{z=-L}^{z=L} + 2ik^2\int_{-L}^{L} \epsilon''(z)|E(z)|^2 dz = 0. \tag{5}$$

Substituting in Eq. (5) the expressions of the electric field at $z = -L$ and $z = L$ respectively we get

$$\mathcal{A} \equiv 1 - \mathcal{T} - \mathcal{R} = \frac{k}{|E_f^-|^2}\int_{-L}^{L} dz|E(z)|^2\epsilon''(z). \tag{6}$$

Furthermore, we assume that $\epsilon''(z)$ is zero everywhere inside the layered structure apart from the interval $0 \leqslant z \leqslant d$ where the defect layer is placed. In this interval it takes a uniform value $\epsilon''(0 \leqslant z \leqslant d) = \epsilon''_{\rm d}(T)$. These simplifications allow us to express the absorption coefficient of Eq. (6) in the form

$$\mathcal{A}(T) = \rho(T)\omega\epsilon''_{\rm d}(T), \tag{7}$$

where $\rho(T) = \mathcal{I}_{\rm d}/|E_f^-|^2$ is the ratio of the integral of light intensity $\mathcal{I}_{\rm d} = \int_0^d dz|E(z)|^2$ at the lossy layer and the incident light intensity. It is obvious from Eq. (7) that $\mathcal{A}(T)$ depends on both the dissipative part of the permittivity and the value of the electric field inside the defect layer. Although the former increases monotonically with the temperature T and thus with the duration time of the incident pulse, this is not true for $\rho(T)$. The latter, which is a unique function of the permittivity, remains approximately constant up to some value of $\epsilon''_{\rm d}$ above which it decreases, leading eventually to a total decrease of the absorption coefficient together with a simultaneous increase of the reflectivity of the structure. This is related to the fact that the increase of $\epsilon''_{\rm d}$ spoils the resonant localized mode [see Fig. 1(c)], which is responsible for high transmittance. Specifically, when the losses due to $\epsilon''_{\rm d}$ overrun the losses due to leakage from the boundaries of the structure, the resonant mode ceases to exist [see Fig. 1(c)] and the structure becomes reflective, i.e., $\mathcal{R} \approx 1$, and $\mathcal{T} \approx 0$ [see Fig. 1(d)]. As a consequence we have that $\mathcal{A} = 1 - \mathcal{T} - \mathcal{R} \approx 0$ and the system does not absorb the high incident energy of the incoming light source but rather reflects it back in space.

In fact, the nonmonotonic shape of the envelope of the scattering field in Fig. 1(c) is a direct consequence of the fact that the structure becomes reflective $\mathcal{R} \approx 1; \mathcal{T} \approx 0$. One has to realize that in the case where both Bragg gratings on the left and right of the defect layer are finite, the field inside each half-space is written as a linear combination of two evanescent contributions with exponentially decreasing and exponentially increasing amplitudes. Their relative weight is determined by the boundary conditions $E(z = -L) = E_0^-(-L)$ and $E(L) = E_0^+(L) = E_f^-\sqrt{\mathcal{T}}$ at the two outer interfaces of the layered structure. In the case of reflective structures these boundary conditions lead to the relation $E(-L) = E_f^- \sim \mathcal{O}(1)$ and $E(L) \approx 0$. It can be shown rigorously that in this case, the field on the left half-space of the structure is dominated originally by the exponentially decaying component while after some turning point z_0 the exponentially increasing component becomes dominant up to the defect layer. After that the field decays exponentially as in the resonant case. Similar scattering field profiles have been found in cases of active (gain) defects [12].

One can use a simple qualitative argument to estimate the condition under which $\mathcal{A}(T)$ continues to increase. As we discuss previously, we assume that the electromagnetic energy losses occur in the lossy defect layer. The dissipated power can be estimated from Eq. (7) to be $\dot{Q} \propto \mathcal{A}|E_f^-|^2 = \omega\epsilon''_{\rm d}\mathcal{I}_{\rm d}$. Due to the energy conservation, the rate of energy dissipation cannot exceed the energy supply provided by the incident wave. The latter is $\mathcal{S}_{\rm in} \propto c|E_f^-|^2$. Taking this constraint into account we get the following upper limit on the field intensity at the defect layer location:

$$\frac{c}{\omega\epsilon''(T)d}|E_f^-|^2 \geqslant |E_{\rm d}|^2. \tag{8}$$

Above we have made the additional approximation that $\mathcal{I}_{\rm d} \sim |E_{\rm d}|^2 d$, where $E_{\rm d}$ is a typical value of the field inside the defect layer.

Next we recall that a resonant mode with a frequency ω inside the bandgap has a Bloch wave number which is imaginary $k = ik''$. The electric field inside the layered structure can be expressed as a pair of evanescent modes, one of which is decaying with the distance z and another one which is growing, i.e., $E(z) = E_f\exp(-k''z) + E_b\exp(k''z)$. To the left of the defect $(-L < z < 0)$, the electric field is dominated by the rising evanescent mode $E(z) \approx E_b\exp(k''z)$ while to the right of the defect $(0 < z < L)$, the dominant contribution is provided by the decaying mode $E(z) \approx E_f\exp(-k''z)$ [13].

The field $E_{\rm d}$ at the location of the defect layer is provided by the rising evanescent mode evaluated at $z = 0$, i.e., $E_{\rm d} \sim E_b$. Therefore, the value of this evanescent mode at the left stack boundary at $z = -L$ is

$$E(-L) \propto E_{\rm d}\exp(-k''L). \tag{9}$$

Comparing Eqs. (8) and (9) we can conclude that if

$$\frac{c}{\omega\epsilon''(T)d}\exp(-2k''L) \ll 1 \tag{10}$$

then the amplitude of rising evanescent mode $E(z = -L)$ at the left stack boundary is much less than amplitude of the incident wave

$$|E(z = -L)|^2 \ll |E_f^-|^2. \tag{11}$$

The latter condition, Eq. (11), implies that the energy density inside the left grating is much smaller than the energy density of the incident wave, and hence only a small portion of the incident light energy $S_I \propto c|E_f^-|^2$ will cross the stack boundary at $z = -L$. In other words, the condition Eq. (10) automatically implies high reflectivity at the stack interface. The condition Eq. (10) for high stack reflectivity (and hence low transmittance and absorption) will always be satisfied if the loss tangent $\epsilon''(T)$ of the defect layer is large enough and/or if the number of layers in the Bragg grating is large enough.

Next, we want to quantify the above arguments. To this end, we calculate explicitly the transport characteristics of our grating structure for an incident laser pulse. Although the analysis can be generalized for any incident pulse shape, in our numerical simulations below, we have assumed for simplicity that the incident laser pulse has a train form [14]

$$\begin{aligned} \mathcal{W}_I(t) &= 0 \quad \text{for} \quad t \leqslant 0 \\ &= w_0 \quad \text{for} \quad 0 \leqslant t \leqslant t_f \\ &= 0 \quad \text{for} \quad t \geqslant 0. \end{aligned} \tag{12}$$

We want to calculate the total energy transmitted, reflected, and absorbed during the duration of the pulse. These can be expressed in terms of the time-dependent transmittance $\mathcal{T}(t)$, reflectance $\mathcal{R}(t)$, and absorption $\mathcal{A}(t)$, which are the main quantities that we analyze below. All other observables can be easily deduced from them. For example, the integrated (over the period of the pulse) absorption $\bar{\mathcal{A}}$ can be defined as

$$\bar{\mathcal{A}} = \frac{\int_{-\infty}^{\infty} dt\mathcal{A}(t)\mathcal{W}_I(t)}{\int_{-\infty}^{\infty} dt\mathcal{W}_I(t)}, \tag{13}$$

while similar expressions can be used for calculating the total (over the period of the pulse) transmittance $\bar{\mathcal{T}}$ and reflectance $\bar{\mathcal{R}}$.

Our starting point is the "rate" equation

$$\frac{d}{dt}T(t) = \frac{1}{C}[\mathcal{A}(T)\mathcal{W}_I(t) + \kappa(T_0 - T)] \tag{14}$$

that describes the heating rate of the defect layer. Above, C is the heat capacity, $\mathcal{W}_I(t) \equiv |\mathcal{E}_I(t)|^2 = |\int d\omega E(\omega)\exp(i\omega t)d\omega|^2$ is the incident light intensity, and κ is the thermal conductance of the defect layer. The first term in Eq. (14) describes the heating process of the lossy layer while the second one corresponds to heat dissipation from the defect layer to the mirror (if any) or to the air. To further simplify our calculations, we assume that the temperature changes are within a domain where both thermal conductance and heat capacity are constants and independent of temperature changes.

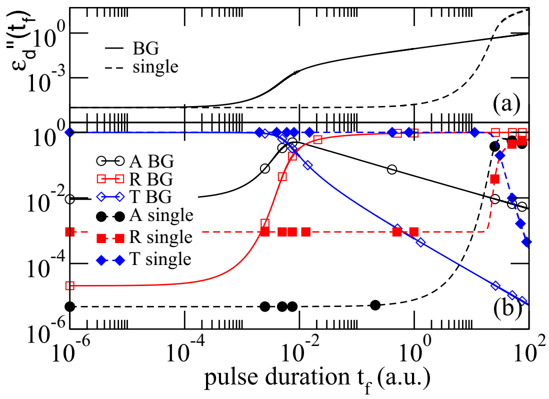

FIG. 2. (Color online) (a) The imaginary part ϵ_d'' of permittivity as a function of pulse duration t_f. The solid line corresponds to the layered structure in Fig. 1, while the dashed line corresponds to the standalone lossy layer. (b) The absorption coefficient $\mathcal{A}_r(t_f)$ (black solid line), reflectance $\mathcal{R}_r(t_f)$ (red [light gray] solid lines), and transmittance $\mathcal{T}_r(t_f)$ (blue [dark gray] solid line) of the layered structure in Fig. 1 vs pulse duration associated with the resonance frequency mode. For longer pulse duration (and larger cumulative energy of the pulse), the absorption $\mathcal{A}_r$ is suppressed and the setup becomes highly reflective ($\mathcal{R}_r \approx 1$). The dashed lines show the respective values (associated with the resonance frequency mode) for the standalone lossy layer, in which case the absorption for pulses with longer duration (and larger cumulative energy) is much higher, while the reflectivity is much lower than those of the layers structured in Fig. 1.

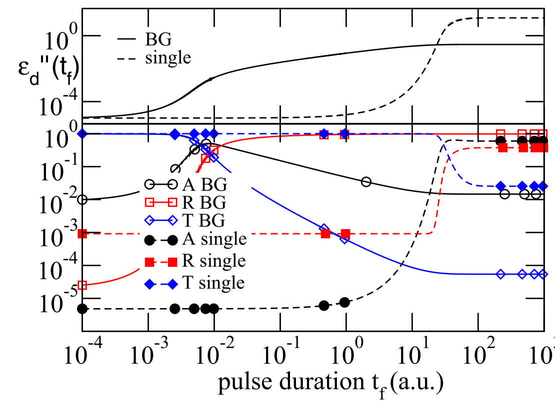

FIG. 3. (Color online) The same as in Fig. 2 but now in the presence of thermal exchange between the defect layer and its surroundings ($\kappa = 0.05$). For longer pulse duration, a steady-state regime is reached, which corresponds to a crossover from energy limiting regime to a power limiting regime.

Substitution of the absorption coefficient from Eq. (7) into Eq. (14) leads us to the following equation:

$$\frac{d}{dt}T(t) = \frac{1}{C}[\omega\varepsilon_{\mathrm{d}}''(T)\rho(T)\mathcal{W}_I(t) + \kappa(T_0 - T)], \qquad (15)$$

which expresses the temporal behavior of the temperature $T(t)$ in terms of the given profile $\mathcal{W}_I(t)$ of the incident pulse. Everything else, e.g., $\epsilon_{\mathrm{d}}''(t)$, $\mathcal{A}(t)$, $\mathcal{T}(t)$, and $\mathcal{R}(t)$, can be directly and explicitly expressed in terms of $T(t)$.

In case that $\kappa = 0$, one can further show that the outcomes can be written in terms of the total incident energy $U_f = \int_0^{t_f} \mathcal{W}_I(t)dt$. Furthermore, using Eq. (15) we get that $T_f = \int_0^{U_f} \mathcal{A}(U)dU/C$. The associated total absorption is $\bar{\mathcal{A}} = (\int_0^{U_f} \mathcal{A}(U)dU)/U_f$, while similar expressions can be derived for the other transport characteristics.

In Fig. 2 we report the outcomes of a direct integration of Eq. (15) for $\kappa = 0$. In this case, the incident thermal energy does not dissipate outside of the defect layer, i.e., the thermal relaxation time is infinite. Therefore, time-cumulative effects are important and thus our structure acts as an energy limiter. In Fig. 2(a) we report the temporal behavior of permittivity ϵ_d'' as a function of the pulse duration t_f. Notice that for train pulses the pulse duration t_f is directly analogous to the total incident energy U_f. We will therefore alternate, in our presentation below, the dependence of ϵ_{d}'', $\mathcal{T},\mathcal{R},\mathcal{A}$ from the pulse duration with the (more natural parameter for an energy limiter) total incident energy of the pulse.

Originally ϵ_d'' is essentially unaffected by the incident energy and the same is true for the resonance mechanism (via the defect mode) that is responsible for high transmittance in the absence of losses. In this domain $\mathcal{T}_r \approx 1$, $\mathcal{R}_r \approx 0$ while there is a slow increase of the absorption $\mathcal{A}_r$, as it can be seen from Fig. 2(b) (solid lines). Once the incident energy (pulse duration time) exceeds some critical value, there is a rather abrupt increase in ϵ_{d}'' which results in the destruction of the resonance mode. Subsequently, the incident energy does not resonate into the structure, leading to a decaying absorption $\mathcal{A}_r \approx 0$, while the same is true for the transmittance $\mathcal{T}_r \approx 0$. At the same time, there is a noticeable growth of the reflectance, which becomes approximately equal to unity $\mathcal{R}_r \approx 1$. For comparison we also plot at the same figure the results of the standalone layer. We find that for large incident energies (pulse durations t_f) the absorption $\mathcal{A}_r(t_f)$ is higher by more than two orders of magnitude as compared to the case of the reflective energy limiter.

We have also performed the same analysis for the case where the thermal conductance κ is different from zero. In Fig. 3 we report the results of the numerical integration of Eq. (15) in the presence of thermal conductivity. For long pulse duration we find a steady-state behavior of the transport characteristics of the reflective energy limiter. The physical nature of the steady-state regime is quite obvious. It corresponds to the situation when the heat released in the defect layer is completely carried away by thermal conductivity. At this point, the temperature of the defect layer stabilizes and the time derivative $dT(t)/dt$ in Eqs. (14) and (15) vanishes. The latter condition determines the steady-state values of the defect layer temperature as a function of the incident light amplitude. In this limiting case our structure acts as a power limiter. For comparison, the results of the standalone lossy layer are also reported in this figure. We find that in the steady-state regime our structure performs superbly, resulting in absorption values which are more than two orders of magnitude smaller than the ones achieved by the standalone lossy layer.

V. CONCLUSIONS

At infrared and optical frequencies, the reflectivity of known uniform materials is well below 90%, especially so when the incident light intensity is dangerously high. So, if we want to build a highly reflective optical limiter, we have to rely on photonic structures which would support some kind of low-intensity resonant transmission via slow or localized modes at photonic bandgap frequencies. If the incident light intensity increases, the respective localized mode must disappear, and the entire photonic structure will behave as a simple Bragg reflector. Here we considered the so-called dissipative mechanism of the localized mode suppression. At first glance, it seems counterintuitive, because the high reflectivity and low absorption are caused by the increase in the loss tangent of the defect layer in Fig. 1. A qualitative explanation for such a phenomenon is that the large value of ϵ'' in the defect layer results in decoupling of the left and the right Bragg reflectors in Fig. 1. Of course, there might be other ways to suppress resonant transmittance when the incident light intensity, or the total energy of the pulse, grow dangerously high. Still, the presented "dissipative" mechanism seems simple and practical.

A key physical requirement to the constitutive materials of the reflective photonic limiter is that the dielectric layers of the Bragg reflectors in Fig. 1 must be lossless and linear. Indeed, if at high-level radiation the Bragg reflector layers also become lossy, the optical limiter will still perform, but it will not be a reflective limiter anymore, because a significant portion of the high-level radiation will be absorbed by the grating. Fortunately, at visible and infrared frequencies there are plenty of available optical materials with negligible losses and nonlinearities that can be used for the construction of the Bragg mirrors.

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[13] We stress that in the case in which both Bragg gratings on the left and right of the defect layer are finite, both evanescent contributions are present in either half-space, although only one if them is dominant on either side. Furthermore, one can show that the presence of both evanescent contributions on either side of the defect layer can provide an energy flux and, hence, a nonzero transmittance.
[14] We have checked that the same qualitative behavior is obtained for other pulse shapes as well.