Non-Linear Localized Modes Give Rise to a Reflective Optical Limiter

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Optical limiters are designed to transmit low intensity light, while blocking the light with excessively high intensity. A typical passive limiter absorbs excessive electromagnetic energy, which can cause its overheating and destruction. We propose the concept of a layered reflective limiter based on resonance transmission via a non-linear localized mode. Such a limiter does not absorb the high level radiation, but rather reflects it back to space. Importantly, the total reflection occurs within a broad frequency range and for an arbitrary direction of incidence. The same concept can be applied to infrared and microwave frequencies.

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The continuing integration of optical devices into modern technology has led to the development of an ever increasing number of novel schemes for efficiently manipulating the amplitude, phase, polarization, or direction of optical beams [1]. Among these manipulations, the ability to control the intensity of light in a predetermined manner is of the utmost importance, with applications ranging from optical communications to optical computing [2, 3] and sensing. As laser technology is making progress, novel protection devices (optical limiters) are needed to protect optical sensors and other components from high-power laser damage [4–6].

Here we focus on the most popular, passive optical limiters. The simplest realization of a passive optical limiter is provided by a single nonlinear layer with the imaginary part $n''$ of the refractive index being dependent on the light intensity $W$. At low intensity, the value $n''(W)$ is relatively small, and the nonlinear layer is transparent. As the light intensity exceeds certain level, the value $n''(W)$ increases dramatically, and the nonlinear protective layer turns opaque. In more sophisticated schemes, the nonlinear layer can be a part of a complicated optical setup. The problem though is that in all cases, the limiter absorbs the excessive power, which might cause overheating or even destruction of the device (a sacrificial limiter). Our goal is, using the existing nonlinear materials, to design a photonic structure which would reflect the excessive power back to space, rather than absorbing it. Such a structure can be referred to as a passive reflective limiter. A free-space realization of a reflective limiter is a layered array reflecting a high intensity light regardless of the direction of incidence and within a broad frequency range.

Our proposal is based on resonance transmission through a nonlinear localized mode. The simplest realization of the above idea is illustrated in Fig. 1 where a nonlinear lossy layer is sandwiched between two linear lossless Bragg reflectors. We will show that if the light intensity is low, the absorption can be small, and the layered structure in Fig. 1 will be transmissive in the vicinity of the localized mode frequency. If the incident light intensity exceeds a certain level, the non-linear layer in Fig. 1 decouples the two Bragg gratings and the entire stack becomes highly reflective – not opaque, as in the case of a stand-alone nonlinear layer. In other words, the high intensity light will be reflected back to space, rather than absorbed by the limiter. Even this simple design provides a broad band protection for an arbitrary direction of incidence. For a given nonlinear material, the intensity limitation of the transmitted light can be controlled by adjusting the layered structure, so that the electromagnetic energy density in the vicinity of the nonlinear layer is either enhanced or attenuated. A problem with the simple design of Fig. 1 is that the low-intensity transmittance only occurs in the vicinity of the localized mode frequency. This problem can be addressed by using more sophisticated photonic structures, for instance, those involving two or more coupled defect layers, as it is done in the case of optical filters [7].

To illustrate our idea, we consider a pair of identical Bragg gratings, each consisting of two alternating layers with real permittivities $\varepsilon_1$ and $\varepsilon_2$, placed in the intervals...
Finally find the amplitudes $E_M$ and $E_{amplitudes}$. Utilizing Eq. (2) together with the optical structure associated with the domain $M$, we have integrated backwards Eq. (1), with the help of a 4th order Runge-Kutta, and obtained the field $E^0(z = 0)$ and its derivative $(dE^0(z)/dz)|_{z = d_1}$, as initial conditions we have integrated backwards Eq. [1], with the amplitudes $E_f^0$ and backward $E_b^0$ propagating amplitudes. Utilizing Eq. (2) together with $M^{(L)}$ we finally find the amplitudes $E_f$ and $E_b$ which allow us to calculate $T$, $R$ and $A$. Note that for a backward map with boundary condition $E^+_f = 1$ we have $|E_f^-|^2 = 1/T$. It is convenient to work with the rescaled variable $\tilde{E}(z) = \sqrt{\gamma}E$. In this representation, Eq. (1) becomes

$$\frac{\partial^2 \tilde{E}(z)}{\partial z^2} + \frac{\omega^2}{c^2} \tilde{E}(z) = 0 \quad (3)$$

where $\tilde{E}(z) = E(z)/\sqrt{\gamma}$, while $\tilde{E}(z) = E(z)/\sqrt{\gamma} = \tilde{E}(z)$, and $\tilde{E}(z) = E(z)/\sqrt{\gamma} = \tilde{E}(z)$. In other words, in this representation, the nonlinear layer has a fixed absorption rate which is equal to unity, the outgoing field boundary associated with the backward map varies as $\tilde{E}_f^- = \sqrt{\gamma}$ when the incident light intensity $W_l$ is $W_l \equiv |E_f^-|^2 = \gamma/T = \gamma|E_f^-|^2$.

In Fig. 2 the effect of the incident intensity $W_l$ on the transmission, reflection and absorption of a resonant localized mode is presented. The Bragg grating used in these simulations consists of 40 layers on each side with alternating permittivities $\epsilon_1 = 4$ and $\epsilon_2 = 9$. The width of the impurity layer $d_\gamma = 1$ and the amplitude $\epsilon$ of the nonlinear permittivity is $\epsilon = 9$. We have confirmed numerically that in the linear case the defect creates a resonant mode at $\omega_r \approx 8.15$ [9] at the band-gap of the grating which is localized around the impurity. We find that as the incident intensity $W_l$ increases (main panel of Fig. 2), the transmittance of this resonant mode decreases, with a simultaneous increase of the absorption. Further increase of $W_l$, results in noticeable growth of the reflectance with a simultaneous decrease of the absorption and transmittance. Eventually both $T$ and $A$
vanishes for moderate values of $W_I$. In other words the system reflects completely the incident radiation. For the sake of comparison we also calculated $T, A$ and $R$ versus $W_I$ for a single non-linear layer with no Bragg reflectors (see inset of Fig. 2). We find that for the same range of moderate values of incident intensity $W_I$, the system rather absorbs the energy instead of reflecting it back to space.

For normal incidence, a further theoretical analysis can be carried out. To this end we assume that the permittivity of the non-linear layer is $\epsilon_r(z) = \epsilon (1 + i\gamma|E(z)|^2)\delta(z)$. This approximation is justified in the case of a thin metallic defect. For the analytical calculation of $T, R$ and $A$, we proceed along the same lines that we have highlighted in the numerical analysis previously. For the sake of generality we will assume that the transport characteristics of the left and right linear subsystems are encoded in the values of their left (right) transmission $t_L(t_R)$ and reflection $r_L(r_R)$ amplitudes. The elements of the transfer matrices $M^{(L)}$ and $M^{(R)}$ (see Eq. (2)) are defined as $M_{11}^{(R/L)} = 1/t_{R/L}^*, M_{12}^{(R/L)} = -r_{L/R}^*/t_{L/R}$, $M_{21}^{(R/L)} = -(r_{L/R}/t_{L/R})^*$, and $M_{22}^{(R/L)} = 1/t_{L/R}$. 

Next we calculate the field amplitudes just before and after the delta defect by utilizing the transfer matrices Eq.(2) associated with the linear segments. For a left incident wave, we have at $z = 0^-$

$$E_f^- = \frac{E_f^- - r_L E_f^+}{t_L^-}; \quad E_b^- = \frac{E_f^+}{t_L^-} - \frac{r_L^* E_f^+}{t_L^-}$$

while at $z = 0^+$ just after the delta defect we have

$$E_f^+ = \frac{t_R E_f^+}{1 - |r_R|^2}; \quad E_b^+ = \frac{t_R^* E_f^+}{1 - |r_R^*|^2}.$$  

Using Eqs. (4) and (5) together with the continuity of the field at $z = 0$ and the suitable discontinuity of its derivative we write the incident and reflected field amplitudes in terms of the transmitted wave amplitude

$$E_f^- = \left\{\frac{1}{\tau_0} - i\left(\frac{1}{\tau_0} - \frac{1}{\tau_f}\right)\gamma|\xi|^2|E_f^+|^2\right\}E_f^+;$$

$$E_f^+ = \left\{\frac{1}{\tau_0} - i\left(\frac{1}{\tau_0} - \frac{1}{\tau_f}\right)\gamma|\xi|^2|E_f^-|^2\right\}E_f^-.$$  

Above, $\tau$ is the transmission amplitude in the absence of the $\delta-$like layer, $\tau_0$ is the transmission amplitude when $\gamma = 0$, and $\xi = t/f + \mu r/|r|^2$. From Eq. (6) we deduce the transmission, reflection and absorption amplitudes. For the transmission and reflection amplitude we get that

$$t = \frac{1}{\tau_0} - i\left(\frac{1}{\tau_0} - \frac{1}{\tau_f}\right)\gamma|\xi|^2|E_f^+|^2; \quad r = \left\{t_L^* - \frac{1}{t_L^*}\right\}(1-r_L).$$

The transmittance, reflectance and absorption can then be calculated as $T = |t|^2$, $R = |r|^2$ and $A = 1 - T - R$. From Eq. (7) we observe that increasing $\gamma$ (we remind that the incident light intensity $W_I \sim \gamma$) results in an increase of the denominator of the transmission amplitude and therefore to a decrease of $T$ (for very large $\gamma$-values it becomes zero). At the same time the reflection amplitude, becomes $r \rightarrow \left(\frac{1}{t_L^*} - \frac{1}{r_L^*}\right)$ corresponding to perfect reflection, i.e. $R \rightarrow 1$. Consequently in this limit we have zero absorption $A = 0$.

Fig. 3 demonstrates the effect of $W_I$ on a resonant localized mode for the case of symmetrically placed Bragg gratings on the left and right side of a $\delta-$like defect. The alternate layers at the Bragg gratings have permittivity $\epsilon_1 = 4$ and $\epsilon_2 = 9$ while the permittivity of the defect layer is $\epsilon = 1.5$. The transport characteristics of the gratings $t_L = t_R$ and $r_L = r_R$ have been calculated numerically and used as inputs in Eqs. (7). We find (see Fig. 3) that the overall behavior of $T, R$ and $A$ is similar to the one observed in the simulations of Fig. 2.

For comparison, we also report (inset of Fig. 3) the behavior of $T, A$ and $R$, for a single non-linear layer (without any Bragg gratings), vs. the incident intensity $W_I$. They are calculated analytically using the continuity of the field and the discontinuity of its derivative at the position of the $\delta-$defect. Specifically,

$$T = \frac{4}{(k\epsilon_0)^2 + (2k\epsilon_0 \gamma |E_f^+|^2)^2}; \quad R = (k\epsilon_0)^2(1 + \gamma^2|E_f^+|^4)/4$$

and $A = k\epsilon_0 \gamma |E_f^+|^2 T$. We find that for moderate $W_I$-values the single non-linear layer is mainly absorptive (inset of Fig. 3) while the structure of Fig. 1 is mainly reflecting the incident light back to space (main panel of Fig. 3).

We have also investigated the efficiency of the proposed limiter in the case of oblique incidence. A representative example in the case of an incident angle $\phi = 60^\circ$ is shown in Fig. 4. The Bragg grating considered in this
dent energy flux \( S \) coming (say from the left) wave that carries an incidence mode which is localized around the impurity layer \( \sim \psi_{+} = \alpha_{+} \exp(-z) + \alpha_{-} \exp(L) \). Let us assume that \( \alpha_{+} \sim \mathcal{O}(1) \). Then the field at the outer boundary of the left grating at \( z = -L \) is \( E_{r}(z = -L) = \alpha_{+} \exp(-L) + \alpha_{-} \exp(L) \sim \alpha_{-} \exp(L) \). At the same time due to continuity at the boundary we expect that the resonance wavefunction must be equal to the incoming field which we assume to take some constant value i.e. \( \alpha_{-} \exp(L) \sim \mathcal{O}(1) \). This can only happen if \( \alpha_{-} \rightarrow 0 \). Finally we recall that the incoming energy flux is given by the Poynting vector \( S \) which in the case of evanescent modes is \( S \sim \psi_{+} \psi_{-} = \alpha_{+} \alpha_{-} \rightarrow 0 \). Therefore there will be no net flux towards the structure and thus \( T = 0 \). Since \( T = 0 \) and \( A = 0 \) we conclude that almost all the incident energy is reflected back i.e. \( R \rightarrow 1 \).

In conclusion, we have examined the scattering problem for a periodic layered structure with an embedded nonlinear defect layer. We presume that the absorption coefficient of the defect layer increases with the light intensity, which is normally the case. We have shown that such a layered structure acts as a self-protecting power limiter. Specifically, at low intensity of the incident light, the entire stack is highly transmissive. When the light intensity increases, the stack transmission decreases. Initially, the fraction of the input power absorbed by the lossy nonlinear layer also increases with the incident light intensity. But when the input power exceeds a certain level, the stack becomes highly reflective within a broad frequency range and regardless of the angle of incidence. In other words, the excessive radiation will be reflected back to space, rather than being absorbed by the limiter, which can prevent overheating and destruction of the limiter. A simplest realization of such a self-protected (reflective) power limiter is provided by a lossy non-linear layer sandwiched between two Bragg gratings. A shortcoming of such a design is that although the high intensity radiation will be reflected back to space within a broad frequency range, the low-intensity transmittance only occurs within a narrow frequency band corresponding to the frequency of the localized mode. This problem can be addressed by using a more sophisticated, structured defect layer, as well as a chain of several coupled nonlinear defects. The latter possibilities are currently under investigation [10].

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[8] The quantities $T, R, A$ are the same for a right incident wave as well, since our structure is reciprocal.

[9] In all numerical simulations in Figs. 2–4 we use units such that $c = 1; d = 1$.

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[11] We assume the incident wave has amplitude $O(1)$ and that due to continuity of the wavefunction at the boundary $E_{res}(z = -L) \sim O(1) \rightarrow E_{res}(z = 0) \sim \exp(L)$

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