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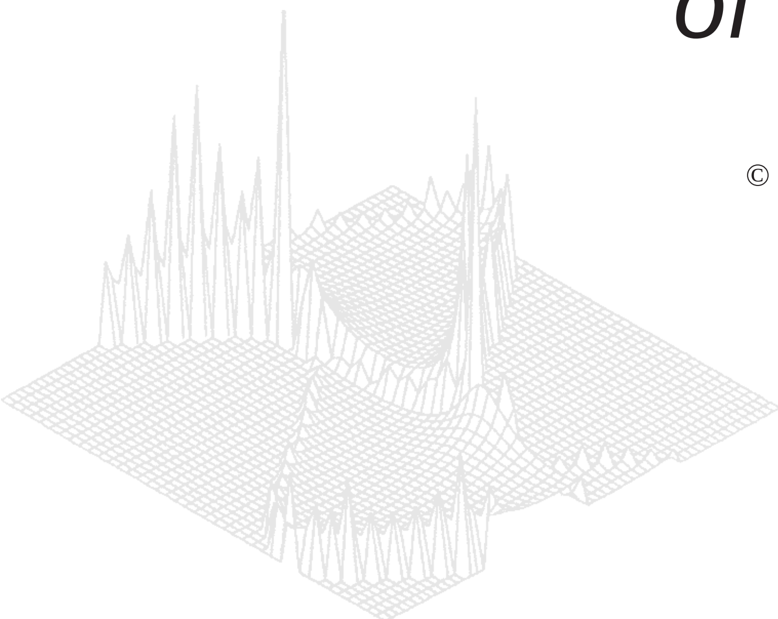

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Light Amplification and Attenuation in Stratified Structures with Complex Refractive Index*

Sergei A. Bulgakov[A,B,C] *and Manuel Nieto-Vesperina*[A]

[A]Instituto de Ciencia de Materiales, CSIC,
Cantoblanco, E-28049, Madrid, Spain.

[B] Present address: School of Physics, University of Sydney,
NSW 2006, Australia.

[C]Permanent address: Radioastronomy Institute,
Ukrainian Academy of Sciences,
4 Krasnoznamennaya Str., 310002 Kharkov, Ukraine.
email: bulgakov@altavista.net

Abstract

We study the electromagnetic wave propagation inside a one-dimensional photonic lattice with dielectric permittivity having a non-zero imaginary part. A detailed investigation of the influence of gains and losses on both the transmission/reflection coefficients and the electromagnetic field distribution is provided. A duality phenomenon between loss and gain is discussed. Upon investigation of the reflected spectrum, we show the existence of huge resonances associated with laser thresholds. Some of these resonances stem from the wave reflection, while others are due to the wave propagation regime. For a regular finite periodic structure with a complex dielectric permittivity, we show a symmetry in the positions of the resonances along the frequency axis.

1. Introduction

Photonic band structures are of great current interest because of their brilliant applications as novel optical high technological materials and because of the possibilities they open to study fundamental physics. It turns out that the main properties of a photonic crystal, the fact that periodic dielectric structures offer the possibility to eliminate propagation of electromagnetic waves through a band of frequencies (the photonic band gap), can solve a major problem of the last decade: to control the optical properties of materials (an extended list of references has been given by Joannopoulos *et al*. 1995). Already, fibre-optics cables, which simply guide light, have revolutionised the telecommunications industry. Lasers, high-speed computers, and spectroscopy, are just a few of the fields next in line to reap the benefits from new optical materials such as *photonic band gap materials*.

Photonic band structures (PBS) are periodic arrays of dielectric materials which cause the dispersion relation for electromagnetic waves propagating in the

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periodic dielectric array to exhibit a band structure. Specifically, frequency band gaps are opened in the dispersion of electromagnetic waves and these gaps inhibit the propagation of electromagnetic waves through the photonic band structure at frequencies in the gaps. The PBS are then technologically useful as frequency filters of electromagnetic radiation. Spontaneous emission mechanisms often are a source of energy loss in lasers and semiconducting devices so that photonic band structures may offer a means to improve the efficiency of such devices.

Photonic crystals are being intensively studied both theoretically and experimentally. It turns out that an intricate overlapping or 'competition' between different scattering mechanisms (Bulgakov and Nieto-Vesperinas 1996) allows one to create artificial dielectric lattices with predictable spectral properties. An important branch of photonic lattice studies is associated with a medium having a complex refractive index. Several works have demonstrated an interesting interplay between the localisation phenomenon and the amplification/attenuation effect which attributes a medium with a nonzero imaginary part of the refractive index. As is known, the localisation of electromagnetic waves in one-dimensional media results from either the random distribution of scatterers or the phase-coherent Bragg resonant conditions. One should expect that, by adding materials with gain or loss into the host dielectric slabs of a stratified medium, the spectral properties and the electromagnetic field distribution inside the structure will be modified. [The spectral and energy properties of both random and periodic lattices made from real dielectics have been widely studied (cf. Brovelli and Keller 1995; Bulgakov and Nieto-Vesperinas 1996; Yariv and Yeh 1977).]

Due to multiple reflections of waves by internal interfaces of the photonic crystal, the interference between counter-propagating waves produces a kind of feedback. In the case of dielectric materials with losses, this feedback will always suppress a wave propagation independently of the kind of medium we use: random or regular, inside or outside a band gap region. However, in the case of an amplifying material, one should expect that the decrease of wave produced by the localisation effect may be compensated.

Most of papers devoted to 1D random lasers (cf. Letokhov 1968) study the probability function of the reflection/transmission coefficient. Asatryan *et al.* (1998) calculated numerically the localisation length with the help of a transfer matrix inside a medium containing randomly distributed impurities with loss and gain. Heinrichs (1997) used the invariant embedding method to study wave propagation inside a random structure containing complex dielectrics. The transmission and reflection coefficients have been studied as functions of the lattice length. Most of works show the existence of a laser threshold and the duality between loss and gain, namely, the gain suppresses the transmission in the same way as the loss does (see Wiersma and Lagendijk 1996; Paasschens *et al.* 1996; Pradhan and Kumar 1994; Zhang 1995). Yariv and Yeh (1977) studied wave propagation through a structure with a periodically alternated imaginary part of the dielectric permittivity. The laser threshold resonances have been observed. However, only a passband region was considered and the switching of losses and gains makes it difficult to understand the roles played by each type of material. Asatryan *et al.* (1998) calculated the poles corresponding to the laser threshold resonances for a wide range of wavelengths; however, due to the disorder in the distribution of impurities with loss or gain, it is difficult to see how the impurity with loss or gain affects separately the wave scattering process. Also, to our knowledge, there is no clear explanation as to why both large gains and losses essentially decrease the transmission. What is the role of multiple scattering in the duality phenomenon? To answer this and other questions above, we use as the most appropriate model a periodic lattice with constant imaginary dielectric background.

2. Model Description and Calculation Procedure

In this paper we address the propagation of an electromagnetic wave through a stratified medium, either periodic or random, with either gains or losses. We demonstrate the existence of new huge threshold resonances in the reflection spectrum different to those previously studied (Yariv and Yeh 1977). We first consider a one-dimensional periodic lattice, and then we address a random distribution of layers.

The length of the structure is L. As for the periodic structure we have $L = Nd$, where N is the number of unit cells of width d (cf. Bulgakov and Nieto-Vesperinas 1996). Each cell contains two dielectric slabs of width d_1 and d_2 ($d_1 = 3d_2$ in our calculations), where $d_1 + d_2 = d$. The dielectric permittivities are $\varepsilon_{1,2} = \varepsilon^r_{1,2} + i\varepsilon^i/\omega^2$, where $\varepsilon^r_{1,2}$ are real and positive ($\varepsilon^r_1 = 1$, $\varepsilon^r_2 = 9$ in our calculations), ω is the frequency of the incident wave, and the term ε^i/ω^2 stands for either gains or losses according to whether $\varepsilon^i < 0$ or $\varepsilon^i > 0$. To justify our model of the dielectric permittivity function, $\varepsilon_{1,2}$, we refer to Haus (1984) who used the gain coefficient $\mathrm{Im}(\varepsilon_{1,2})$ as a function of frequency ω such as $1/\{1 + [(\omega - \omega_0)/\omega_g]^2\}$, where ω_0 and ω_g are the centre frequency of the gain medium and its width, respectively. Most other works on media with random losses and gains consider that ε^i is independent of the frequency. In our present study, we have considered a dependence of $\mathrm{Im}(\varepsilon_{1,2}) = \varepsilon^i/\omega^2$. This allows us to study characteristics of the wave scattering process inside higher frequency bands. We shall deal with a linear polarisation of the electromagnetic field such that the electric vector is parallel to the slab interfaces. We do not restrict the application of our results by considering any concrete dielectric material. Information about realistic active and passive dielectrics can be found in Haus (1984), see also Wiershma and Lagendijk (1996).

To calculate both the transmission and reflection coefficients, and the field distribution inside the lattice layers, we shall apply the transfer matrix in the form given by Born and Wolf (1975). The numerical algorithm used was described in detail by Bulgakov and Nieto-Vesperinas (1998).

Fig. 1 shows the transmittance, $\log T$, calculated for different lattice lengths versus ε^i at a frequency either inside a passband (Fig. 1) or inside a band gap (Fig. 1, inset). (In this paper we use values of the dimensionless frequency $W = \omega d/c$ where c is the velocity of light.) In Fig. 1 each curve has a maximum at $\varepsilon^i = \varepsilon^i_{max} < 0$, where ε^i_{max} depends on the lattice length $L = Nd$. On varying L, the maximum value of T also changes. In Fig. 1 we observe that the lattice of 10 unit cells yields a larger transmission peak than any other structure in the figure. This absolute maximum of the T-value in Fig. 1 corresponds to $\varepsilon^i_{max} \approx -2{\cdot}6$. However, another local maximum in the transmission value associated with a lattice of $L = 140d$ also appears at $\varepsilon^i \approx -0{\cdot}1$.

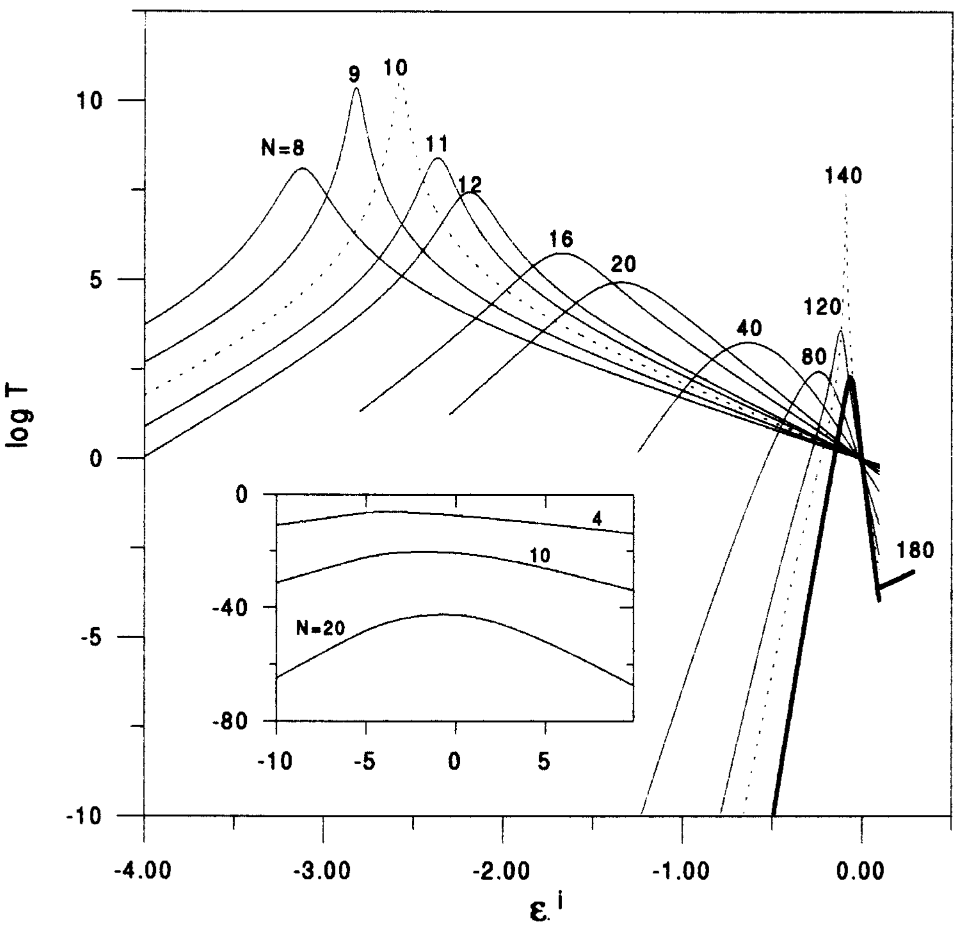


Fig. 1. Transmission coefficient T as a function of ε^i, calculated for different lattice lengths $L = Nd$ at both the passband frequency $W = 4{\cdot}2$ and the band gap frequency $W = 2{\cdot}1$ (inset).

In addition, we observe in Fig. 1 the common cross point of all curves at $\log T = 0, \varepsilon^i = 0$. This means that, for a given particular configuration, the lattice is transparent ($T = 1$) for real dielectric permittivity ($\varepsilon^i = 0$) at this particular frequency $W = 4{\cdot}2$. In the case of losses ($\varepsilon^i > 0$), one has $\log T < 0$. At a fixed $\varepsilon^i > 0$, the transmission is greater the shorter the lattice is.

The transmittivity curves calculated at the frequency inside a band gap (see the inset in Fig. 1 with $W = 2{\cdot}1$) are quite different to those inside a passband (see Fig. 1). Maximum transmittivity is attained at some value $\varepsilon^i_{max} < 0$. This maximum of $\log T$ shifts increasingly from $\varepsilon^i = 0$ the shorter the lattice is; for example, $\varepsilon^i_{max} \approx -4, -2{\cdot}5, -0{\cdot}9$ is associated with lattices $N = 4, 10, 20$ respectively). Also, at fixed ε^i, the longer lattice possesses the lower transmission value.

A remarkable property of the transmission coefficient as a function of ε^i, both inside the passband (Fig. 1) and band gap (Fig. 1 inset) is the following: any departure of ε^i far from the value ε^i_{max} results in a decrease of the transmission coefficient independently of the sign of ε^i. Thus, both large gains and losses suppress wave transmission. As a result, at large enough $|\varepsilon^i|$, one cannot distinguish whether a medium possesses losses or gains by just examining its transmission spectrum: Specifically, in our investigation, we have observed that $T(-\varepsilon^i) \approx T(\varepsilon^i)$, $\varepsilon^i \to +\infty$.

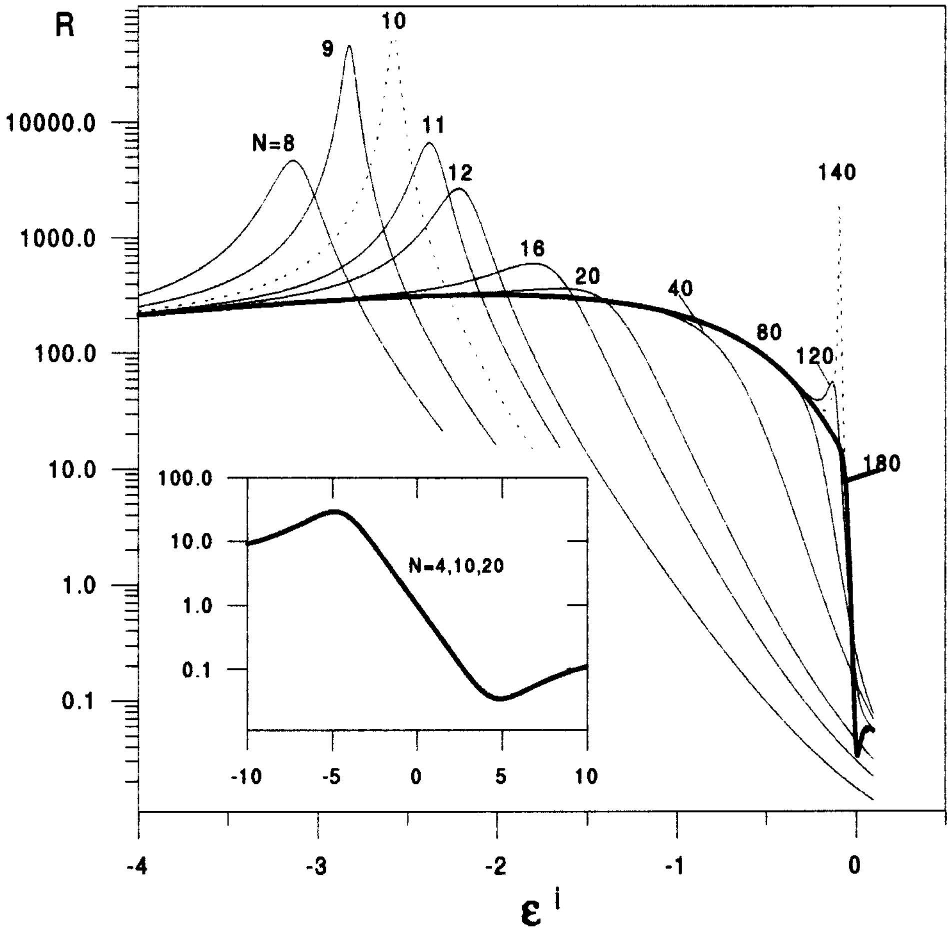


Fig. 2. Reflection coefficient R as a function of ε^i calculated for different lattice lengths $L = Nd$ at both the passband frequency $W = 4{\cdot}2$ and the band gap frequency $W = 2{\cdot}1$ (inset).

Concerning the reflection coefficient, Fig. 2 shows its variation versus ε^i. Calculations are done at the same frequency as for the transmission T of Fig. 1. Inside a passband (Fig. 2) the positions of $\log R$ maxima on the ε^i-axis, for N = 8, 9, 10, 11, 12, 16, 120, 140, coincide with those of the maxima of the corresponding lattices in Fig. 1. We observed in our investigation that for large $|\varepsilon^i|$, $R(\varepsilon^i)$, associated with different lattice lengths, asymptotically tend to R =constant (this is not visible in Fig. 2 because of the scale). The value of this constant depends on the frequency and sign of ε^i: In contrast with the behaviour of the transmission coefficient in Fig. 1, R is always $R > 1$ for a medium with gains ($\varepsilon^i < 0$), and $R < 1$ if $\varepsilon^i > 0$ (losses).

3. Discussion of the Results

Before discussing further the results of Figs 1 and 2, we show in Fig. 3 the dependence of $\log R$ on both the dimensionless frequency $W = \omega d/c$ and ε^i for a lattice of length $L = 20d$. In the vicinity of the plane $\varepsilon^i = 0$, one observes a sequence of regions where $\log R$ oscillates with W. They correspond to passbands. These regions are separated by 'smooth valleys' which are associated with band gaps ($R = 1$). At any fixed value of W, the reflection coefficients saturate as $|\varepsilon^i| \to \infty$.

Another very important characteristic of the reflection coefficient in Fig. 3 is the existence of huge peaks ($R \to \infty$) at $\varepsilon^i < 0$ and far from the plane $\varepsilon^i = 0$; and there are deep minima ($R \to 0$) in the region $\varepsilon^i > 0$ at positions symmetric to those peaks with respect to $\varepsilon^i = 0$.

Although not shown here, we have also studied the surface $\log T(W, \varepsilon^i)$. In contrast to Fig. 3, $T \to 0$ as $|\varepsilon^i| \to \infty$ for any value of the frequency (cf. also Fig. 1), and no pronounced dips are observed in the region $\varepsilon^i > 0$. However, resonances with $T \gg 1$ were recognisable in the region $\varepsilon^i < 0$ at exactly the same coordinates (W, ε^i) at which $R \to \infty$ in Fig. 3.

Let us now consider a lattice with disorder in the distribution of scatterers. We apply a random removal of layers according to the procedure previously used (Bulgakov and Nieto-Vesperinas 1996, 1998), namely, we randomly substitute about 60% of slabs with $\varepsilon_2 = \varepsilon_2^r + i\varepsilon^i/\omega^2$ by those with $\varepsilon_1 = \varepsilon_1^r$. Such random substitutions make the structure periodic on average (other statistical properties of this random procedure have also been studied by Bulgakov and Nieto-Vesperinas 1996). Fig. 4 shows the function $\log R(W, \varepsilon^i)$, as in Fig. 3, for this random structure. Averages have been performed over 30 realisations of the structure. Further averaging very slowly smooths this result, but at a very large computing cost. The spectral characteristics of the random lattice of real dielectric permittivity (i.e. $\varepsilon^i = 0$) were discussed by Bulgakov and Nieto-Vesperinas (1996, 1998).

Next, we proceed to discuss the properties of the reflection and transmission coefficients given in Figs 1–4. In the vicinity of $\varepsilon^i = 0$, say, $|\varepsilon^i| \ll W^2\varepsilon_{1,2}^r$, the phase coherence yielding Bragg resonances is still preserved. Therefore, in a passband region, the phase relations help wave transmission. However, since the forward and backward waves simultaneously increase or decrease as $\exp(-\varepsilon^i|z|)$, depending on the sign of ε^i, then, in this region, the wave propagates throughout the lattice while its amplitude exponentially decreases for a medium with losses. This explains why $T < 1$ decreases with L in Fig. 1 for $\varepsilon^i > 0$.

Concerning the reflection coefficient inside a passband (cf. Fig. 2, $\varepsilon^i > 0$), although always $R < 1$ for a medium with losses, in general, the reflection decreases slower than T as ε^i departs from 0, and saturates as $\varepsilon^i \to +\infty$ [we shall later discuss the cause of minima in the dependence $R(\varepsilon^i)$ for $\varepsilon^i > 0$ (Fig. 2) which are also visible as deep holes in Fig. 3]. The explanation is as follows: as ε^i increases, it is mainly the waves reflected from the first few interfaces which contribute to the reflected power. The reflection that occurs deep in the lattice progressively decreases because of absorption as ε^i increases. Therefore,

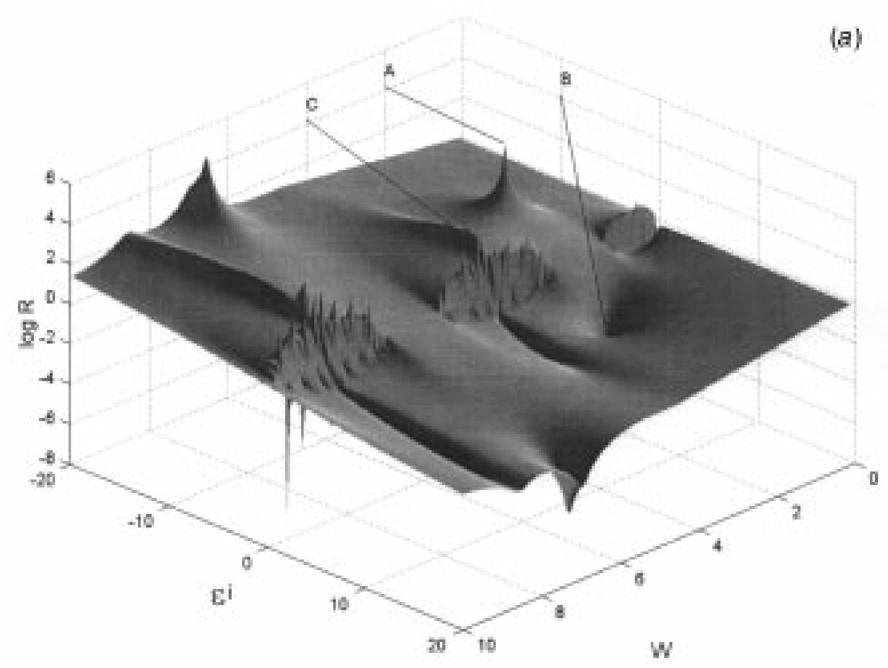
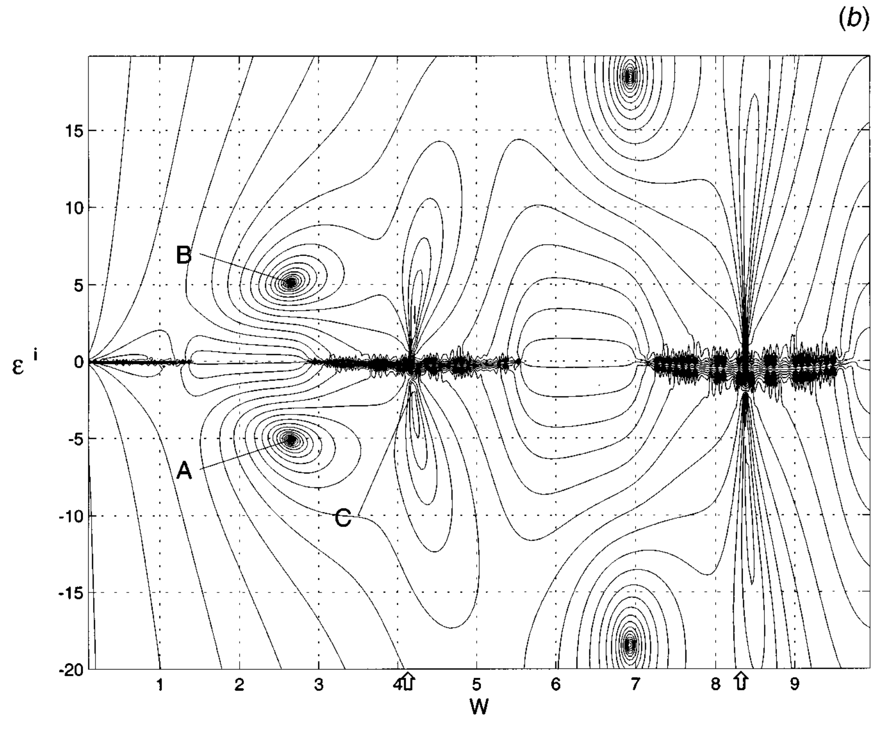


Fig. 3. Logarithm of the reflection coefficient as a function of both the frequency W and ε^i, calculated for the lattice of 20 periods: (a) a 3D view and (b) a contour plot.



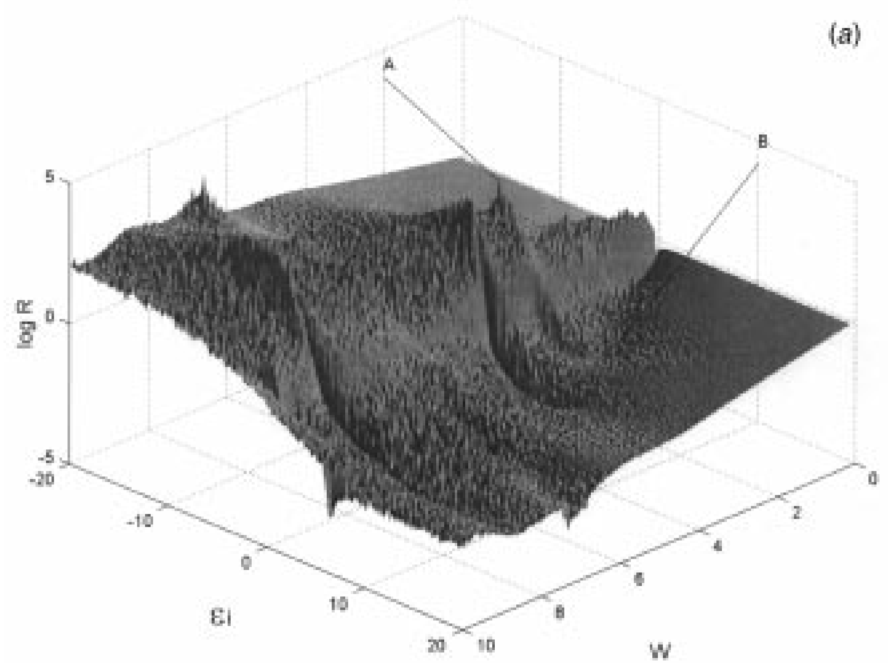
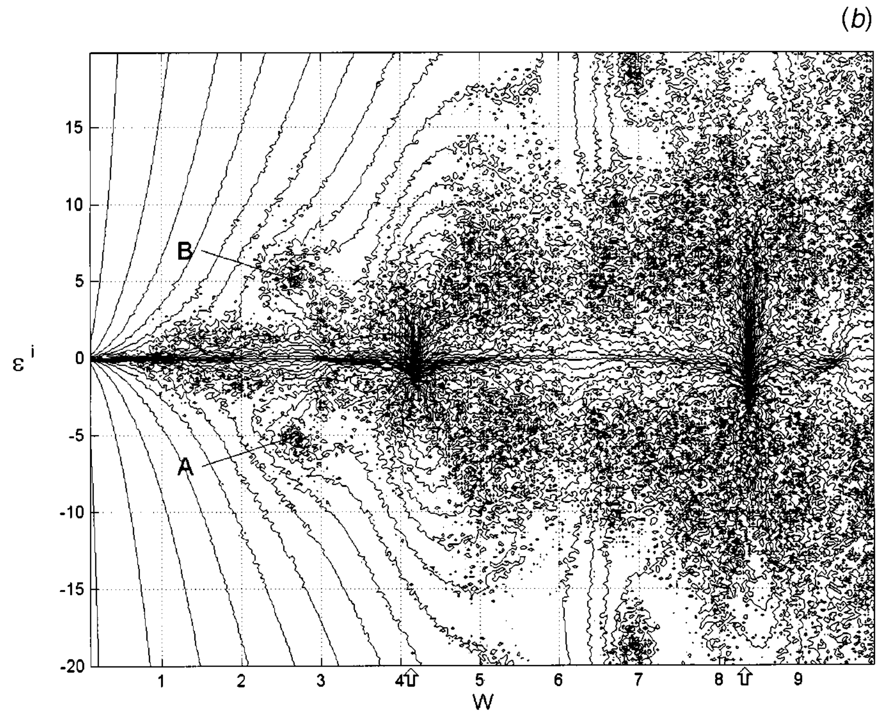


Fig. 4. Same as Fig. 3, but for the structure with disorder in the distribution of scatterers. The length of the structure is $20d$: (a) a 3D view and (b) a contour plot.



the reflection coefficient saturates at the same value, independently of the lattice length.

On the other hand, for a medium with gains, at a passband frequency (cf. Figs 1 and 2 for $\varepsilon^i < 0$), a complicated behaviour of the transmission and reflection coefficients as functions of ε^i results from the fact that at a certain value $\varepsilon^i = \varepsilon^i_{max}$ a *laser threshold* is reached (see Kogelnik and Shank 1972; Letokhov 1968; Yariv and Yeh 1977), so that the energy that escapes from the lattice is compensated by the field amplitude amplification described by the exponential factor $\exp(|\varepsilon^i z|)$.

At the laser threshold we observe both huge reflection and transmission resonances. Any further increase of $|\varepsilon^i|$, $\varepsilon^i < 0$, results in a situation in which strong reflection, due to gain material, does not allow the wave to propagate inside the lattice. Thus, the transmission continuously decreases with an increase of $|\varepsilon^i|$ far from the laser threshold. In contrast, if $\varepsilon^i \to -\infty$, the reflection coefficient saturates at a value that does not depend on the lattice length (see Fig. 2 inset) due to the contribution to the reflected power from the first few layers at the lattice entrance.

We also see in Figs 1 and 2 that the threshold value ε^i_{max}, at a fixed frequency inside a pass band, decreases with an increase of the lattice length L. Notice that the reflection coefficient behaves in the same way as the transmission coefficient (compare Figs 1 and 2 for $\varepsilon^i < 0$). Specifically, the reflectivity is due to reflection throughout the lattice. For a short lattice, a larger gain coefficient is required in order to compensate the energy that escapes from the lattice by transmission and reflection.

The peaks of the transmission values $T_{max} = T(\varepsilon^i_{max})$ depend non-monotonically on the lattice length. In Fig. 1 one observes two local maxima T_{max} of $\log T$ versus L at the frequency $W = 4{\cdot}2$. To explain such behaviour, let us recall that inside a pass band of a lossless finite lattice ($\varepsilon^i = 0$) there is a set of spectral peaks whose magnitude and, therefore, their width depend on the number of unit cells N of the lattice. The peaks are equidistantly distributed along the frequency axis inside the pass band. On varying N, one does not shift the pass band position (Bulgakov and Nieto-Vesperinas 1996), however, the positions of the peaks inside the pass band will change. Although the frequency $W = 4{\cdot}2$ lies inside the second pass band, the transmissions associated with lattices of different length are different: It is evident that for lattices with $L = 10d$ and $L = 140d$, the value $W = 4{\cdot}2$ coincides with the corresponding resonant peak inside the pass band, whereas for other values of L, it does not.

The following question is now in order: how do T_{max} and R_{max} behave as functions of the lattice length, providing that $T = 1$ at $\varepsilon^i = 0$ for a lattice of any length? As has been shown by Bulgakov and Nieto-Vesperinas (1996), at the frequency of a single scatterer resonance (SSR) (the 1D analogue of a Mie resonance), namely, at $W = m\pi d/d_2\sqrt{\varepsilon^r_2}$, m being an integer, and $T = 1$, $R = 0$ independently of the lattice length. Our calculations show that for $W = 4{\cdot}18879$, associated with a single scatterer resonance of $m = 1$, both T_{max} and R_{max} tend monotonically to infinity as the lattice length L increases, and the abscissa of these maxima, ε^i_{max}, tends to zero with $L \to \infty$. For the lattice of $N = 20$ unit cells (Fig. 3), the SSR (marked by arrows in Fig. 3*b*) manifest themselves as either high 'ridges' or deep 'valleys' according to whether $\varepsilon^i < 0$ or $\varepsilon^i > 0$.

Concerning the aforementioned huge resonant peaks and dips of $R(\varepsilon^i, W)$, symmetrically placed with respect to $\varepsilon^i = 0$ (Figs 3*a* and 3*b*), their coordinates ε^i and W turn out to be independent of the lattice length. Specifically, we observe that the minima positions of the reflection coefficient remain exactly the same whatever the lattice length. However, the maxima positions slightly change as the lattice has $N = 1, 2$ or 3 unit cells, but they quickly stabilise for $N \geq 5$. Therefore, by analysing the expressions for the transmission and reflection coefficients of the structure that contains only one unit cell, i.e. two dielectric slabs ε_1 and ε_2 embedded in a host medium of real dielectric constant, $\varepsilon_{1,2} = \varepsilon^r_{1,2} + i\varepsilon^i$, we shall derive an estimate for the positions of those maxima and minima. On using the transfer matrix method, as applied by Born and Wolf (1975), we find the following transcendental equations:

$$|m_{11} + m_{12} - m_{21} - m_{22}| = 0, \text{ for minima of } R;$$

$$|m_{11} + m_{12} + m_{21} + m_{22}| = 0, \text{ for maxima of both } T \text{ and } R;$$

Here m_{ij}, $i, j = 1, 2$ are the transfer matrix elements as shown in Born and Wolf (1975, see p. 69). As mentioned above, the coordinates (ε^i, W) of minima that satisfy the first of the two equations above are valid for any value of N. The pairs (ε^i, W) that solve the second equation are approximately the coordinates of maxima of both R and T if $N > 1$. However, we have observed that for the lattice with $N \geq 10$, the positions of maxima and minima of the reflection coefficient are completely symmetric with respect to $\varepsilon^i = 0$. On the other hand, it can be shown analytically that, in contrast to R, the transmission coefficient possesses no dips.

We have observed, however, no correlation between the positions of pass bands and band gaps (recognisable at $\varepsilon^i = 0$ in Fig. 3) and those of 'peaks' and 'dips'. The distance between these resonances varies from one pair to another. The maxima and minima of R occur far from $\varepsilon^i = 0$, so one can assume that, due to large values of $|\varepsilon^i|$, the system looses any phase coherence such as that which gives rise to Bragg resonances. In contrast to what occurs in the vicinity of $\varepsilon^i = 0$, the origin of maxima of R and the independence of their positions of the lattice length, stems from the wave reflection from the first few layers.

To discuss a cause of the minima of R (or 'dip') observable in Fig. 3 for $\varepsilon^i > 0$, one must recall the 'negative feedback' that stems from the interaction between counter propagating waves in a medium with losses. The latter means that the energy, returned to the system by feedback and being of opposite sign, exactly compensates the energy that enters in the lattice.

To better understand the nature of the 'peak' and 'dip' resonances, Fig. 5 shows the field amplitude $|E|$ distributions associated with the points A and B in Fig. 3 (thin and thick curves in Fig. 5 respectively). The lattice length is $L = 20d$. Observe that both curves are practically identical and may be transformed into each other by a simple shift along the $|E|$-axis. Most importantly, the field behaves like in a band gap frequency, as it exponentially decreases inside the lattice (cf. Bulgakov and Nieto-Vesperinas 1998).

If one plots the field amplitude distribution for the other kind of laser threshold resonance, namely, the one associated with the wave propagation at a pass band

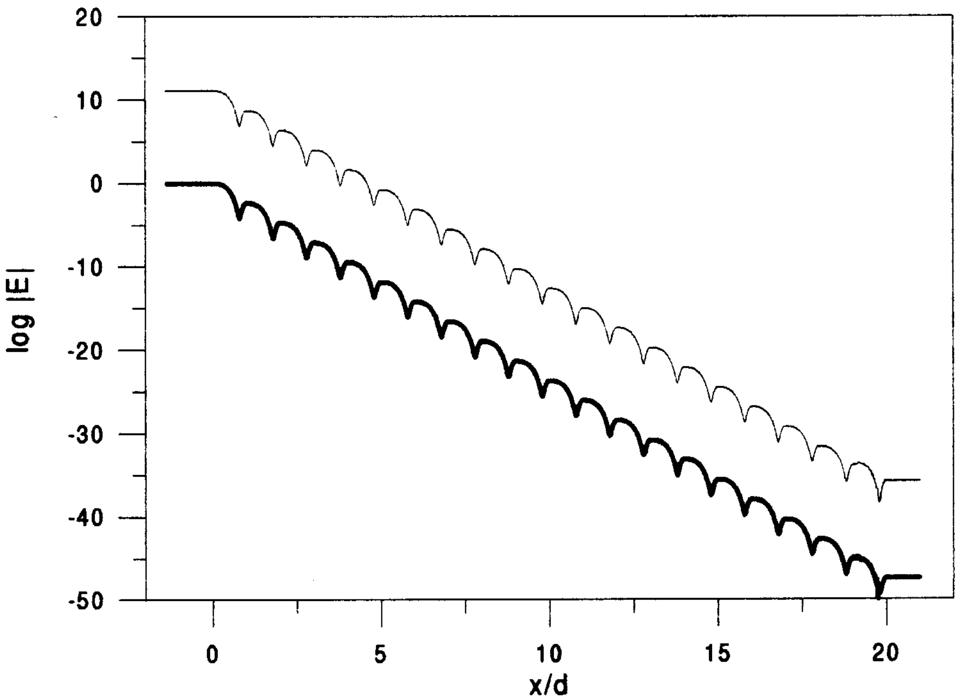


Fig. 5. Field distributions inside the lattice of $L = 20d$ associated with the following different points marked in Fig. 3: the 'peak resonance' (letter A in Fig. 3)—thin curve; and the 'dip' (letter B in Fig. 3)—thick curve.

frequency (cf. point C in Fig. 3), one gets (not shown here) a curve similar to that already reported by Yariv and Yeh (1977, see Fig. 7).

Concerning Fig. 4 for the random stratified medium with complex dielectric constant, one observes that, on average, the disorder does not destroy the existence of SSR in the vicinity of $\varepsilon^i = 0$ (compare with Fig. 3*a* where one such resonance is marked by the letter C). However, the SSR (arrows in Fig. 4*b*) are swamped by the disorder as $|\varepsilon^i|$ increases. This results from the known fact that the disorder causes wave localisation even at the pass band frequency (cf. Bulgakov and Nieto-Vesperinas 1998). In contrast, all huge resonances and deep minima shown in Fig. 3 for the regular structure (cf. points A and B respectively) are clearly manifested in Fig. 4 at the same positions as in Fig. 3, despite the randomness. Since this kind of laser threshold is due to wave interaction with the first few boundaries at the entrance of the layered structure, it does not matter how the introduced disorder shuffles the rest of the layers deeper in the structure.

4. Conclusion

We have shown in this paper that neither randomness in the distribution of scatterers nor periodicity of 1D photonic structures are principal causes of the duality phenomenon between loss and gain. According to our investigation, numerous internal interfaces in a stratified structure stipulate the existence of counter propagating electromagnetic waves. The duality phenomenon results from the mutual contribution of these waves into the final field distribution along the structure. Moreover, it can be shown that even a structure consisting of two interfaces only (e.g. a Fabry–Perot resonator) and containing dielectric material with either gain or loss is an appropriate model for observation of the duality phenomenon.

Two kinds of laser threshold resonances have been reported. One is due to a wave propagation regime through the whole structure. This kind of resonance can be observed in the parameter region where phase coherence helps a wave to propagate. Another laser threshold is due to wave interaction with a few first structure interfaces. The positions of these last resonances are almost independent of the lattice length and the strength disorder in the distribution of scatterers.

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References

Asatryan, A. A., Nicorovici, N. A., Botten, L. C., de Sterke, C. M., Robinson, P. A., and McPhedran, R. C. (1998). *Phys. Rev.* B **57**, 13535.

Born, M., and Wolf, E. (1975). 'Principles of Optics' (Pergamon: New York).

Brovelli, L. R., and Keller, U. (1995). *Opt. Commun.* **116**, 343.

Bulgakov, S., and Nieto-Vesperinas, M. (1996). *J. Opt. Soc. Am.* A **13**, 500.

Bulgakov, S., and Nieto-Vesperinas, M. (1998). *J. Opt. Soc. Am.* A **15**, 503.

Haus, H. (1984). 'Waves and Fields in Optoelectronics', Chapter 9 (Prentice–Hall: Englewood Cliffs, NJ).

Heinrichs, J. (1997). *Phys. Rev.* B **56**, 8674.

Joannopoulos, J. D., Meade, R. D., and Winn, J. N. (1995). 'Photonic Crystals: Moulding the Flow of Light' (Princeton University Press).

Kogelnik, H., and Shank, C. V. (1972). *J. Appl. Phys.* **43**, 2327.

Letokhov, V. S. (1968). *Sov. Phys. JETP* **26**, 835.

Paasschens, J. C. J., Misirpashaev, T. Sh., and Beenakker, C. W. J. (1996). *Phys. Rev.* B **17**, 11887.

Pradhan, P., and Kumar, N. (1994). *Phys. Rev.* B **50**, 9644.

Wiersma, D. S., and Lagendijk, Ad. (1996). *Phys. Rev.* B **54**, 4256.

Yariv, A., and Yeh, P. (1977). *J. Opt. Soc. Am.* **67**, 438.

Zhang, Z.-Q. (1995). *Phys. Rev.* B **52**, 7960.

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