Waveguides of defect chains in photonic crystals

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We develop a two-stage multiple-scattering formalism for the calculation of photonic bands generated by chains of defects in photonic crystals consisting of nonoverlapping spheres, and for the calculation of the transmission of light through straight and bent waveguides of such chains. We apply the method to a specific example which demonstrates that transmission through a bent waveguide occurs with the same efficiency as for a straight waveguide.

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I. INTRODUCTION

The most desirable property of photonic crystals is the presence of an absolute (omnidirectional) photonic band gap which in turn promises important applications in optoelectronics.[1–4] One possible application of photonic crystals, namely, the effective and low loss guiding of photons in integrated optical devices, has also received considerable attention in recent years. So far, most theoretical and experimental efforts have been focused on two-dimensional (2D) photonic-crystal-based waveguides[5–15] but work on three-dimensional (3D) structures has also been reported.[16–20] One way of waveguiding in photonic crystals is that of a coupled-defect waveguide: a single defect gives rise to a resonance state of the electromagnetic (EM) field, at a frequency within the frequency gap of the crystal, localized about the defect. In a periodic arrangement of defects (defect chain), the interaction of neighboring defects creates a narrow band of states (defect band) around the resonance frequency of the single defect, allowing propagation of the field along the chain. Due to the nature of this band, it can be studied by means of a tight-binding (TB) description.[19,20]

From the computational methods available for studying 3D photonic crystals of nonoverlapping spherical scatterers, such as colloidal crystals, opals, metallodielectric crystals, etc., those based on the Korringa-Kohn-Rostoker (KKR) method are the most accurate and computationally efficient, whether in the traditional[21,22] or in the on-shell formulation.[23,24] Extensions of the KKR method to deal with photonic crystals containing single defects[25] and randomly distributed defects within the coherent-potential approximation[26,27] have also been formulated.

In this paper we present a two-stage multiple-scattering KKR method for the treatment of single defects, clusters of defects, and defect waveguides in photonic crystals consisting of nonoverlapping spheres in a homogeneous host medium. We first calculate the propagator functions of the periodic system (reference system). If the reference system under consideration does not support states of the EM field over a range of frequencies, as is the case of a photonic crystal for frequencies within an absolute band gap, the propagators decay exponentially with the distance and can be calculated directly, without the need of elaborate Ewald-type transformations. In the second stage we calculate the propagator functions of the system containing the defects, again directly in real space, as only a small number of nearest neighbors suffice to give good convergence of the lattice sums involved. The restriction of the multiple-scattering equations to only a few neighbors gives a TB character to the method. We note that, due to the rapid decay of the propagator functions, most of the matrices involved in the computation are sparse; exploiting this we have been able to develop an order-N method for the study of photonic crystals with defects (the computer time scales linearly with the number N of the inequivalent defects). In Sec. II we describe the method with the emphasis placed on the problem of waveguiding. In Sec. III we present numerical results for the defect modes and the transmission properties of a given waveguide.

II. THEORY

A. Multiple-scattering equations

We consider a harmonic, monochromatic EM wave of angular frequency ω which is incident on a sphere embedded in a homogeneous medium of different dielectric function, ϵ. The sphere is centered at the origin of coordinates. The electric-field component of the EM wave has the form $\mathbf{E}(\mathbf{r};t)=\mathrm{Re}[\mathbf{E}(\mathbf{r})\exp(-i\omega t)]$; outside the sphere, $\mathbf{E}(\mathbf{r})$ is expanded in spherical waves as follows (see, e.g., Refs. 23 and 28):

$$\mathbf{E}(\mathbf{r})=\sum_{l=1}^{\infty}\sum_{m=-l}^{l}\left[\frac{i}{q}a_{lm}^{0E}\nabla\times j_l(qr)\mathbf{X}_{lm}(\hat{\mathbf{r}})+a_{lm}^{0H}j_l(qr)\mathbf{X}_{lm}(\hat{\mathbf{r}})+\frac{i}{q}a_{lm}^{+E}\nabla\times h_l^{+}(qr)\mathbf{X}_{lm}(\hat{\mathbf{r}})+a_{lm}^{+H}h_l^{+}(qr)\mathbf{X}_{lm}(\hat{\mathbf{r}})\right], \quad (1)$$

where $q=\sqrt{\epsilon}\omega/c$, c being the velocity of light in vacuum; j_l and h_l^+ are the spherical Bessel and Hankel functions, respectively; and $\mathbf{X}_{lm}(\hat{\mathbf{r}})$ are vector spherical harmonics. The associated magnetic-field component of the EM wave can be obtained through Maxwell equations and we need not write it down explicitly. The first two terms of Eq. (1) represent the incident (incoming) wave on the sphere and the last two

represent the scattered (outgoing) wave. In the remainder of the paper we will use the composite index L to describe collectively the indices lmP with $P=E,H$. By applying the proper boundary conditions at the surface of the sphere, we obtain a relation between the expansion coefficients of the incident and the scattered fields, as follows,

$$a_L^+ = T_L\, a_L^0, \tag{2}$$

where T_L are the elements of a scattering matrix (**T** matrix). Explicit forms for the T_L of a single sphere can be found elsewhere.[28]

Let us now consider an assembly of nonoverlapping spheres centered at sites $\mathbf{R}_i$ in a homogeneous host medium. An outgoing wave about $\mathbf{R}_{i'}$ [the wave field is described by terms similar to the last two of Eq. (1) with coefficients $b_L^{+i'}$] can be written as an incoming wave expanded about $\mathbf{R}_i$ [the wave field is described by terms similar to the first two of Eq. (1) with coefficients $b'^{i}_{L}(i')$] by means of the equation

$$b'^{i}_{L}(i') = \sum_{L'} \Omega^{ii'}_{LL'}\, b_{L'}^{+i'}. \tag{3}$$

Explicit expressions for the so-called free-space (which here means a homogeneous host medium characterized by a real dielectric function) propagator functions, $\Omega^{ii'}_{LL'}$, for the EM field can be found in Ref. 26. We note that, by definition, $\Omega^{ii'}_{LL'}$ equals zero for $i=i'$.

We now introduce the propagator functions $D^{ii'}_{LL'}$ for the assembly of spheres in the host medium, which give the coefficients [in an expansion such as that of Eq. (1)] of the wave incident on the sphere at $\mathbf{R}_i$, due to an outgoing wave from the sphere at $\mathbf{R}_{i'}$, noting that an outgoing wave from the i'th sphere can reach the ith sphere directly, or indirectly, after scattering any number of times by any number of spheres (including those at $\mathbf{R}_i$ and $\mathbf{R}_{i'}$). One can easily prove by iteration the following equation:

$$D^{ii'}_{LL'} = \Omega^{ii'}_{LL'} + \sum_{i'',L''} \Omega^{ii''}_{LL''} T^{i''}_{L''} D^{i''i'}_{L''L'}, \tag{4}$$

where T^i_L are the elements of the **T** matrix for the sphere at $\mathbf{R}_i$. For a periodic arrangement of spheres, one usually calculates the Fourier transform of $D^{ii'}_{LL'}$, denoted by $D_{LL'}(\mathbf{k})$ ($\mathbf{k}$ is the Bloch wave vector), and integrates the latter over the Brillouin zone (BZ) to obtain $D^{ii'}_{LL'}$. However, unless the spheres are absorbing, the numerical calculation over the BZ requires a very dense mesh of $\mathbf{k}$ points due to singularities in $D_{LL'}(\mathbf{k})$.[27,29,30] On the other hand, as a rule, the calculation of $D^{ii'}_{LL'}$ on the basis of Eq. (4) involves a summation over a large number of lattice sites because the free-space propagator functions, $\Omega^{ii'}_{LL'}$, decay slowly with the distance $|\mathbf{R}_i - \mathbf{R}_{i'}|$. But there are exceptions to this rule: the lattice sum of Eq. (4) can be truncated to a relatively small number of nearest neighbors around the site $\mathbf{R}_i$ for frequencies within a region over which EM waves do not propagate in the host medium. This is true, for example, in the case of a metal host at frequencies below its bulk plasma frequency, in which case $\Omega^{ii'}_{LL'}$ decay exponentially with the distance.[31] The lattice sum in Eq. (4) may also be rapidly convergent in the case of a photonic crystal which possesses an absolute frequency gap because in this case the propagator functions $D^{ii'}_{LL'}$ decay exponentially with distance at frequencies which lie within the absolute gap (see Sec. III).

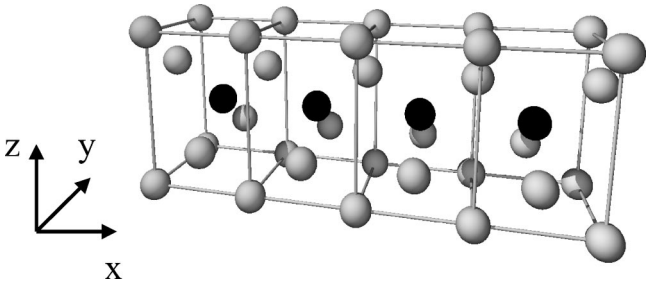

FIG. 1. A diamond structure of spheres. A linear chain of defect spheres (black spheres) is introduced along the [110] direction which is taken as the x axis.

If we know the propagator functions $D^{r\,ii'}_{LL'}$ for an arbitrary system of nonoverlapping spheres in a homogeneous host medium (reference system) we can obtain $D^{ii'}_{LL'}$ for a perturbed system (a number of spheres of the reference system are replaced with spheres of different size and/or dielectric function) by solving the following equation:

$$D^{ii'}_{LL'} = D^{r\,ii'}_{LL'} + \sum_{i'',L''} D^{r\,ii''}_{LL''} \Delta T^{i''}_{L''} D^{i''i'}_{L''L'}, \tag{5}$$

where ΔT^i_L is the difference between T^i_L of the perturbed system and T^{ri}_L of the reference system. Obviously, the lattice sum of Eq. (5) is restricted to the sites where spheres of the reference system have been substituted. In the absence of propagating states in the reference system, $D^{r\,ii'}_{LL'}$ decay exponentially with distance and, therefore, the lattice sum of Eq. (5) can be truncated to a few nearest neighbors around the site $\mathbf{R}_i$.

B. A defect chain

We consider a 3D photonic crystal consisting of nonoverlapping spheres centered at $\mathbf{R}_i = \mathbf{R}_n + \mathbf{t}_\alpha$, where $\mathbf{R}_n$ is a (Bravais) lattice vector and $\mathbf{t}_\alpha$ a nonprimitive vector within the unit cell of the crystal; we may have more than one sphere in the unit cell in which case we have a corresponding number of $\mathbf{t}_\alpha$ vectors. In what follows $i \equiv (n,\alpha)$. We assume that the photonic crystal possesses an absolute frequency gap. By substituting an infinite number of spheres along a specific crystallographic direction of this crystal (host crystal), with spheres of different radius and/or dielectric function (but the same for all defects), we construct a periodic linear chain of defects inside the crystal (see Fig. 1). Evidently all defect spheres are situated at the same position $\mathbf{t}_\alpha$ within the unit cell of the crystal, which we denote as $\alpha \equiv 0$. The above may be seen as a system with one-dimensional (1D) periodicity (along the defect chain) with the defect spheres at $R_n = nR$

where R denotes the period of the chain. The propagator functions of the perturbed crystal shown in Fig. 1 can be obtained from those of the host crystal, considered as the reference system, through Eq. (5). Putting $i=(n,0)\equiv n$ by which we denote the nth sphere of the chain, and similarly $i'=(n',0)\equiv n'$, we have

$$D_{LL'}^{nn'}=D_{LL'}^{r\,nn'}+\sum_{n'',L''} D_{LL''}^{r\,nn''}\Delta T_{L''}D_{L''L'}^{n''n'}, \tag{6}$$

where ΔT_L is the difference between T_L of the defect spheres and T_L^r of the host ones; evidently the sum over n'' includes the spheres of the chain only.

Taking advantage of the 1D periodicity of the system described by Eq. (6), we multiply Eq. (6) by $\exp[-ikR(n-n')]$ and take the sum over n', to obtain

$$D_{LL'}(k)=D_{LL'}^{r}(k)+\sum_{L''} D_{LL''}^{r}(k)\Delta T_{L''}D_{L''L'}(k), \tag{7}$$

where

$$D_{LL'}^{r}(k)=\sum_{n'} \exp[-ikR(n-n')]D_{LL'}^{r\,n\,n'} \tag{8}$$

with a similar expression for $D_{LL'}(k)$. We rewrite Eq. (7) as

$$\sum_{L''} [\delta_{LL''}-D_{LL''}^{r}(k)\Delta T_{L''}]D_{L''L'}(k)=D_{LL'}^{r}(k). \tag{9}$$

The poles of $D_{LL'}(k)$, for given k, determine the eigenfrequencies of the EM field associated with the chain for this k. And since $D_{LL'}^{r}(k)$ has no poles over the frequency gap of the unperturbed crystal, these eigenfrequencies are given by the roots of

$$\det[\delta_{LL'}-D_{LL'}^{r}(k)\Delta T_{L'}]=0 \tag{10}$$

for every k in the 1D BZ: $-\pi/R<k\leqslant\pi/R$.

Let us now consider a finite chain consisting of N defect spheres, obtained by the replacement of spheres of the host crystal; the defect chain need not be linear in the general case (see below). If we set $M_{LL'}^{ii'}=\delta_{LL'}\delta_{ii'}-D_{LL'}^{r\,ii'}\Delta T_{L'}^{i'}$, we can write Eq. (5) as

$$D_{LL'}^{ii'}=\sum_{i'',L''} (M^{-1})_{LL''}^{ii''}D_{L''L'}^{r\,i''i'}. \tag{11}$$

Since $D_{LL'}^{r\,ii'}$ decay exponentially with the distance $|\mathbf{R}_i-\mathbf{R}_{i'}|$, we drop all terms of Eq. (11) which involve interactions beyond first neighbors. Then the propagator which connects the first and the last spheres of the defect chain becomes

$$D_{LL'}^{N1}=\sum_{L''} [(M^{-1})_{LL''}^{N1}D_{L''L'}^{r\,11}+(M^{-1})_{LL''}^{N2}D_{L''L'}^{r\,21}]. \tag{12}$$

We are particularly interested in $D_{LL'}^{N1}$, which is taken as a measure of the transmission of an EM wave through the defect chain. We need not invert the matrix $M_{LL'}^{ii'}$ as a whole in order to calculate $D_{LL'}^{N1}$ from Eq. (12). Since we consider only first-neighbor interaction in the defect chain, the **M** matrix has a block-tridiagonal form: only the diagonal ($i=i'$) and the neighboring upper and lower diagonal blocks ($i'=i\pm1$) have nonzero elements. In this case, $(M^{-1})_{LL'}^{N1}$, $(M^{-1})_{LL'}^{N2}$, can be calculated in such a way that the time required for their computation scales linearly with the number N of defect spheres in the chain.[32]

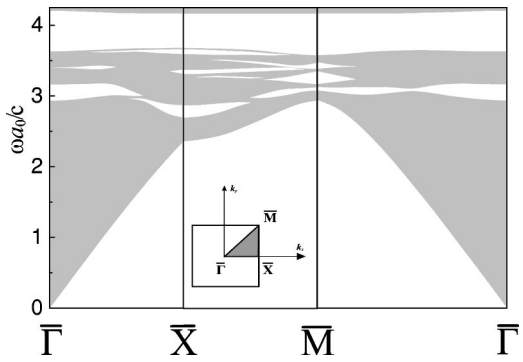

FIG. 2. Projection of the frequency band structure, of a crystal consisting of dielectric spheres ($\epsilon_s=12.96, S/a_0=0.25$) arranged as in a diamond structure in air ($\epsilon=1$), on the SBZ of the (001) surface (shown in the inset).

From $D_{LL'}^{N1}$ we obtain the matrix elements of the scattering path operator

$$\tau_{LL'}^{N1}=T_L D_{LL'}^{N1}T_{L'}, \tag{13}$$

which gives the L component of the scattered (output) EM wave from the last (Nth) sphere due to an incident (input) L' wave on the first defect sphere.

III. APPLICATIONS

We consider a 3D photonic crystal consisting of nonoverlapping spheres of dielectric constant $\epsilon_s=12.96$ in air ($\epsilon=1$). The spheres are arranged as in a diamond crystal of lattice constant a. We view the crystal as a stack of layers parallel to the xy plane. The periodicity of the layers parallel to this plane is described by a 2D square lattice defined by the primitive vectors $\mathbf{a}_1=a_0(1,0,0)$ and $\mathbf{a}_2=a_0(0,1,0)$, where $a_0=a\sqrt{2}/2$ is the distance between second nearest neighbors in the diamond structure. A basis of two spheres of radius $S=a_0/4$, at (0,0,0) and $a_0(1/2,0,\sqrt{2}/4)$, defines the two planes of spheres of a layer. The $(n+1)$th layer along the z axis is obtained from the nth layer by a simple translation described by the primitive vector $\mathbf{a}_3=a_0(1/2,1/2,\sqrt{2}/2)$ (see Fig. 1).

Figure 2 shows the projection of the frequency band structure of the EM field in this crystal on the surface Brillouin Zone (SBZ) of the (001) surface, calculated using the computer program of Ref. 28. The shaded regions extend over the frequency bands of the EM field: at any one frequency within the shaded region, for a given $\mathbf{k}_\parallel$ [the component of the reduced wave vector within the SBZ of the (001) surface], there exists at least one propagating mode of the

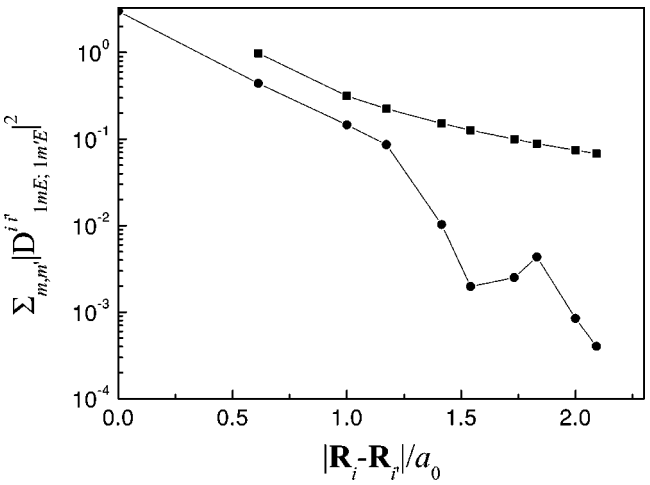

FIG. 3. The circles show $\Sigma_{m,m'}|D^{ii'}_{1mE;1m'E}|^2$ as a function of the distance $|\mathbf{R}_i-\mathbf{R}_{i'}|$, for the crystal of Fig. 2. The squares show $\Sigma_{m,m'}|\Omega^{ii'}_{1mE;1m'E}|^2$. Both curves were obtained for $\omega a_0/c=3.92$ (middle of the absolute band gap shown in Fig. 2).

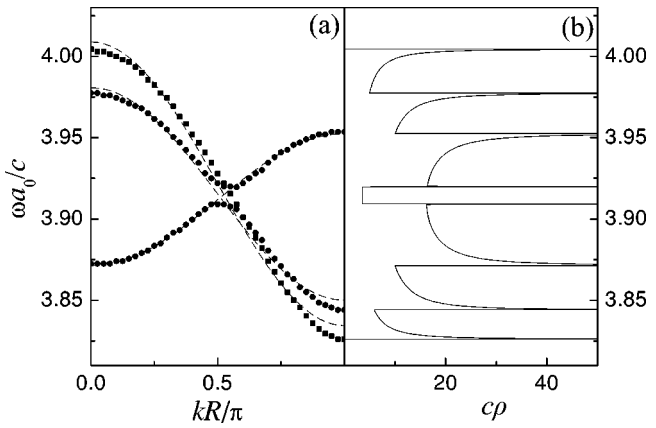

FIG. 4. (a) Defect bands generated by an infinitely long linear chain of defect spheres ($\epsilon_s=12.96, S/a_0=0.21$) along the [110] direction of the crystal described in Fig. 2. The squares/circles are bands of A/B symmetry, respectively, calculated from Eq. (10). The dashed lines are fits of the (unhybridized) bands to Eq. (14). (b) Density of the EM modes corresponding to the defect bands.

EM field in the infinite crystal. The blank areas correspond to frequency gaps. One clearly sees that we obtain an absolute frequency gap extending from $\omega a_0/c=3.68$ to $\omega a_0/c=4.16$. We verified that this is indeed so by calculating the band structure at a sufficient number of $\mathbf{k}_\parallel$ points within the SBZ. Over this frequency region, there are no propagating states of the EM field. At frequencies near the middle of the absolute frequency gap, the propagator functions $D^{r\,ii'}_{LL'}$ obtained from Eq. (4) decay rapidly to zero as the distance $|\mathbf{R}_i-\mathbf{R}_{i'}|$ increases, in contrast to the free-space propagator functions $\Omega^{ii'}_{LL'}$ which decrease very slowly with the distance (see Fig. 3). Therefore, the procedure described in Sec. II applies.

We introduce, to begin with, a linear defect chain, by replacing N spheres lying along the [110] direction by others of different radius: $S_d=0.21a_0$ (see Fig. 1); for an infinitely long chain, $N\to\infty$. We note that the distance $R=a_0$ between first-neighbor spheres along the defect chain corresponds to the second-neighbor distance in the diamond structure and that the distance between second neighbors in the linear chain corresponds to the distance between eighth neighbors in the diamond structure; therefore, at frequencies near the middle of the absolute frequency gap where $D^{r\,ii'}_{LL'}$ decrease rapidly with the distance, considering first-neighbor interaction between the defect spheres in the second-stage scattering [Eq. (6)] should be a realistic approximation to the problem under consideration. The same applies also to the bent waveguide described below.

In Fig. 4(a) we show the frequency bands associated with the linear infinite chain (we refer to them as defect bands), calculated from Eq. (10). The calculation was done for a sufficient number of k points in the region $0\leqslant k\leqslant \pi/R$; the bands for $k<0$ are of course symmetric to the ones shown. It can be seen that the defect bands extend over a narrow frequency region about the center of the absolute gap, which ensures that they are reliably calculated by the procedure described in Sec. II. We note that all three bands, which are nondegenerate, result from modes localized on the spheres of the chain interacting weakly between them. We confirmed that the corresponding eigenmodes are made up, almost entirely, of electric-dipole terms [E terms with $l=1$ and $m=-1,0,+1$ in Eq. (1)] and magnetic dipole terms [H terms with $l=1$ and $m=-1,0,+1$ in Eq. (1)]. In this respect the bands resemble the electric-dipole bands associated with a chain of spheres in a homogeneous medium with a negative dielectric constant, studied in Ref. 19, except that in the present case the reduced symmetry of the system does not allow any degeneracy in the bands. Here, the symmetry of the defect bands is determined by the C_{1h} point group, and the corresponding propagating modes (Bloch waves) of the EM field belong to either A or B, the two 1D irreducible representations of this group.[33] In fact, we obtain three nondegenerate dipole bands of which one has the A symmetry and the other two have the B symmetry. All three bands disperse with the wave number k in the manner characteristic of TB modes, and are described very well by [dashed lines in Fig. 4(a)]

$$\omega_\nu(k)=\overline{\omega}_\nu+W_\nu\cos(kR)\qquad \nu=1,2,3, \tag{14}$$

where $\overline{\omega}_\nu$ and W_ν are constants which we determined by fitting them to the exact results as shown in Fig. 4(a). We expect a degree of hybridization between defect bands of the same (B) symmetry, and this naturally leads to the exact bands shown by the squares and circles in Fig. 4(a). It is also worth noting that in the present case the modes are mostly of the H type rather than the E type met in the photonic insulator of Ref. 19. The density of EM modes resulting from these bands is easily obtained from

$$\rho(\omega)=\frac{1}{\pi}\sum_\nu\left[\frac{d\omega_\nu(k)}{dk}\right]^{-1}, \tag{15}$$

where ν runs over the different defect bands [see Fig. 4(b)]. We remember that $\rho(\omega)d\omega$ gives the number of EM modes per unit length of the chain between ω and $\omega+d\omega$. It is worth noting that $\rho(\omega)$ diverges at those frequencies where the slope of a band vanishes, although $\rho(\omega)$ is integrable (Van Hove singularities), a fact which might lead to a measurable Lamb shift in atomic transitions of approximately the same frequency when these transitions occur within the photonic crystal.[34]

Finally, we describe the transmission of light across a chain of finite length (N spheres) through the following quantity (we may call it the $L'\equiv l'm'P'$ propagation coefficient):

$$t^{N1}(L')\equiv\sum_{L}|\tau^{N1}_{LL'}|^2, \tag{16}$$

which constitutes a measure of the transmission of light from the first sphere of the chain through the chain to the last (Nth) sphere of it. $t^{N1}(L')$ is not to be confused with the transmission coefficient as ordinarily defined: the ratio of the energy flux outgoing towards infinity from the Nth sphere to the energy flux incident on the first sphere coming from some source outside the chain of spheres. This, the usual transmission coefficient, will depend on the way light comes to or is generated at the first sphere and on the way light is taken away from the Nth sphere, and in this paper we shall not consider this problem. It is perhaps easier to explain the physical meaning of $t^{N1}(L')$ as follows. We first note that, for the finite chain of N spheres we are considering, the wave field is a standing wave field, because over the given frequency region light cannot propagate in the host photonic crystal. Now, $T_L D^{N1}_{LL'}$ tells us that if about the first sphere, *locally*, there is an outgoing wave (L') of unit amplitude, then the *locally* outgoing field about the Nth sphere will have a component L with amplitude $T_L D^{N1}_{LL'}$. Multiplying this quantity with $T_{L'}$ [see Eq. (13)] is a formal way of relating the locally outgoing L' wave from the first sphere with a fictitious incident L' wave on this sphere. What $\tau^{N1}_{LL'}$, given by Eq. (13) (or for that matter $D^{N1}_{LL'}$), tells us is how the field at the end of the chain relates to the field at the beginning of the chain. Summing over the available $L=lmP$ scattering channels at the Nth sphere to obtain $t^{N1}(L')$ [see Eq. (16)] provides a simpler overall measure of the propagation (transmission) of light along the waveguide from the first to the Nth sphere. It is evident that unlike the ordinary transmission coefficient this quantity can be greater than unity. It is a useful quantity. By comparing $t^{N1}(L')$ of a bent waveguide with that of a straight waveguide (as in Fig. 5), we can tell whether the bend hinders the transmission of light. It is also very useful in the study of Anderson localization due to disorder: if the properties of the spheres along a chain vary randomly, then localization due to disorder will trap the light and $t^{N1}(L')$ will vanish exponentially with N. We shall not be concerned with this problem here. A treatment of disorder and of localization along these lines has been done for the corresponding electronic problem (transport of electrons along a chain of atoms) by a number of authors (see, e.g., Ref. 35).

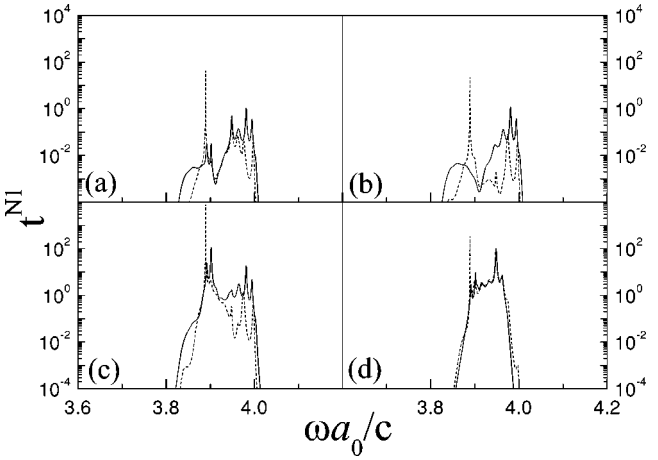

FIG. 5. Transmission along a straight (solid lines) and a bent (broken lines) waveguide for an incoming spherical wave corresponding to (a) $P'=E$, $l'=1$, and $m'=1$; (b) $P'=E$, $l'=1$, and $m'=0$; (c) $P'=H$, $l'=1$, and $m'=1$; and (d) $P'=H$, $l'=1$, and $m'=0$.

In Fig. 5 the solid lines give $t^{N1}(L')$ for a chain of ten defect spheres along a straight line, along the [110] direction, as shown in Fig. 1. The broken lines give $t^{N1}(L')$ for a chain consisting of two sections, a straight line of six defect spheres along the [110] direction and a second straight line of six defect spheres along the $[\bar{1}10]$ direction, the two meeting at an angle of 90° (11 defect spheres in total: one sphere is common to both legs). The two top diagrams refer to an incident E wave ($P'=E$). In Fig. 5(a) $l'=1$, $m'=1$; in Fig. 5(b) $l'=1$, $m'=0$. The results for an E wave with $l'=1$, $m'=-1$ are practically identical to those of Fig. 5(a). The two bottom diagrams refer to an incident H wave. In Fig. 5(c) $l'=1$, $m'=1$; in Fig. 5(d) $l'=1$, $m'=0$. The results for an H wave with $l'=1$, $m'=-1$ are practically identical to those of Fig. 5(c). We observe that in the case of an H wave with $l'=1$, $m'=0$ the propagation coefficient, as defined by Eq. (16), is the same for the straight and bent waveguides at all frequencies. In all other cases the variation of the propagation coefficient with the frequencies is different for the two waveguides, although the difference between the two is on average rather small. It is worth emphasizing that this result has been obtained using the exact EM field and it does not depend on the parametrized representation of the bands shown in Fig. 4(a).

IV. CONCLUSION

We have developed a two-stage multiple-scattering method which allows an exact calculation of waveguiding (transmission of light) across a chain of defects in a realistic photonic crystal consisting of nonoverlapping spheres. The defect chain introduces narrow bands in the absolute gap of the pure crystal. We have shown that an EM signal can propagate through a bent waveguide with more or less the same efficiency as it does through a straight waveguide.

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