Photon and acoustic phonon coupling in phoxonic crystals

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ABSTRACT

We consider phoxonic crystals in relation with their application to photon-phonon interaction. The main interests of such structures are the possibility to confine simultaneously optical and elastic waves in cavities and waveguides and to engineer the photonic and phononic dispersion relations of waveguides. A variety of coupling mechanisms are discussed for exaltation or quenching, including classical photo-elastic coupling as a volume interaction effect, and couplings introduced by moving boundaries or resonator dimensions.

Keywords: Photonic crystal, phononic crystal, optomechanics, Brillouin scattering, electrostriction, elastooptic effect

1. INTRODUCTION

Brillouin and acousto-optic interactions have traditionally been considered under a weak coupling regime for photons and acoustic phonons. The photo-elastic tensor is classically involved and the effect can be described by the interaction of three waves that satisfy both phase-matching and energy conservation. Two photons and one phonon are then involved. The phonon frequency is much smaller than that of the two photons, but its wavelength is commensurate. More recently, the optomechanical effect has been shown to exist in cavities that confine simultaneously photons and phonons, providing an additional coupling mechanism.

The efficiency of opto-acoustic effects depends directly on the local energy density of the involved waves. Hence there are two obvious directions in which to look for enhanced interactions: spatial confinement (for cavities and waveguides) and low group velocities (for waveguides). Spatial confinement has to be simultaneous for light and sound to be effective. The modal distributions of optical and acoustic waves should also be matched so that the interaction is optimized, a condition which depends both on the geometry of the nanostructure and on the choice of materials. Low group velocities are useful in order to increase the interaction time for a given interaction length.

Phoxonic crystals have the unique potential to provide simultaneous confinement and dispersion control of photons and phonons. This fascinating concept has appeared only recently.[1] The central x in the neologism phoxonic stands for t and n at once, meaning that a phoxonic crystal is simultaneously a photonic and a phononic crystal.[2] In the literature, the term optomechanical crystal is also employed.[3] The condition that is generally looked for is that a complete photonic and a complete phononic band gap are present simultaneously. The choice of materials and structure is of utmost importance in order to find an adequate phoxonic crystal configuration, and this quest has been the subject of many papers recently. We review in this paper the properties of the different phoxonic crystal structures that have been proposed, including 2D and 3D crystals, 1D phoxonic crystal strips, and arrays of holes, pillars or resonators of the membrane type. Prior to this presentation, we discuss some possible mechanisms for photon-phonon interaction in dielectric nanostructures.

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2. INTERACTION OF PHOTONS WITH ACOUSTIC PHONONS

2.1 Types of interaction

Photons and acoustic phonons – or optical and elastic waves – interact inside matter in a rather indirect fashion. From a classical point of view, they are waves propagating in a certain medium, e.g., a solid dielectric. Optical waves are often modeled via Maxwell equations as electromagnetic waves, with the macroscopic electric and magnetic fields considered as the basic physical quantities. From a microscopic point of view, however, photons interact with the local electronic distribution, at the atomic scale. The corresponding macroscopic picture is that of the induced dielectric polarization and magnetization of materials. Similarly, elastic waves satisfy the dynamical equations of continuum mechanics, limited to purely elastic terms for small strains. A microscopic view of the same system leads to phonons as solutions of lattice dynamics equations. Acoustic phonons, with frequencies well below any internal molecular resonance and with wavelengths much longer than inter-atomic distances, can be thought as equivalent to continuum elastic waves in a certain limiting process. It remains that elastic waves or acoustic phonons can both be viewed as collective vibrations of the atomic lattice, with displacements measuring the departure from equilibrium positions. The mass of atoms is mostly located in nuclei, but the electronic distribution in the lattice is obviously strongly involved in the inter-atomic potentials that enter lattice dynamics equations – or in classical terms the stiffness or elastic constants of the medium.

This gross physical picture set, the point we wish to make is that the role of electrons is indirect though essential in the coupling of optical and elastic waves inside matter. A propagating elastic wave causes the electronic distribution to be modulated on an acoustic wavelength scale and at acoustic frequencies, and hence the polarization of the medium. This coupling mechanism is the essence of the photo-elastic effect, Brillouin light scattering and electrostriction of phonons from photons, as we discuss in Section 2.2.

A type of photon-phonon interaction different from the photo-elastic effect is provided by optomechanical coupling. As its name suggests, this type of interaction was introduced originally for the coupling of the localized optical modes with the mechanical modes of a resonator. It has a connection with radiation pressure, as we underline in Section 2.3.

An effect related to optomechanical coupling is the modulation of reflected and transmitted optical waves at a moving boundary between two media. If the vibration of the boundary is caused by elastic waves, then this offers a further possible type of acousto-optical interaction, as we discuss in Section 2.4.

2.2 Photo-elastic effect and electrostriction

The photo-elastic effect as used, e.g., in acousto-optics, is a phenomenological approach well adapted to experiments.[4,5] Suppose an elastic wave propagates in a medium, with the vector displacement field written as

$$u_i(\mathbf{r},t) = u_i(\mathbf{r}) \exp(\imath(\Omega t - \mathbf{K}\cdot\mathbf{r})) + \text{c.c.} \tag{1}$$

If the displacements remain small, according to the Pockels effect they will induce a change in the inverse of the permittivity tensor at optical frequencies that is linear with the excitation (first-order approximation)

$$(\Delta\epsilon_r^{-1})_{ij} = p_{ijkl} u_{k,l} \tag{2}$$

with p_{ijkl} the rank-4 photo-elastic tensor, resulting in the nonlinear polarization

$$P_i = \epsilon_0 \chi_{ijkl} E_j u_{k,l} \tag{3}$$

with

$$\chi_{ijkl} = -(\epsilon_r)_{im}(\epsilon_r)_{jn} p_{mnkl} \tag{4}$$

In all equations of this paper, we generally use Einstein's convention of implied summation on repeated indices.

The above formulation is especially useful in locally homogeneous media, i.e., if the elastic and optical properties do not change on the scale of the displacements, while at the same time they can depend on space and time. Indeed, the driven optical wave equation can be taken as

$$\nabla\times(\nabla\times\mathbf{E}) + \frac{1}{c^2}\frac{\partial^2(\epsilon_r\mathbf{E})}{\partial t^2} = -\mu_0\frac{\partial^2\mathbf{P}}{\partial t^2} \tag{5}$$

where the part of the permittivity that is directly influenced by elastic waves, in the right-hand side of the equation, has been purposely separated from the part that does not depend on it, in the left-hand side. In the frame of acousto-optics, where it is generally assumed that the elastic wave is unperturbed by optical waves, this non-linear equation can be transformed into a set of coupled linear equations, using an expansion in diffracted optical waves, and solved by matrix algebra.[6]

A slightly different view, however, is to consider the elasto-optical interaction as a three-wave interaction, where the elastic wave can evolve upon interaction with optical waves, and where for simplicity only one diffracted optical wave is considered. This picture can be equivalently described as the interaction of two photons and a phonon, as in the Brillouin effect, making the introduction of phase-matching easier.[7,8] Specifically, we represent the optical wave as the superposition

$$\mathbf{E}(\mathbf{r},t) = \mathbf{E}^{(1)}(\mathbf{r})\exp(\imath(\omega_1 t - \mathbf{k}_1\cdot\mathbf{r})) + \mathbf{E}^{(2)}(\mathbf{r})\exp(\imath(\omega_2 t - \mathbf{k}_2\cdot\mathbf{r})) + \text{c.c.} \tag{6}$$

Phase-matching in the three-wave interaction is strictly achieved if $\Omega = \omega_1 - \omega_2$ and $\mathbf{K} = \mathbf{k}_1 - \mathbf{k}_2$. Equation (5) leads to

$$(\nabla\times(\nabla\times\mathbf{E}^{(2)}))_i + \frac{\omega_2^2}{c^2}\epsilon_r E_i^{(2)} = -\frac{(\omega_1-\Omega)^2}{c^2}\chi_{ijkl}E_j^{(1)}u_{k,l}^* \tag{7}$$

and

$$(\nabla\times(\nabla\times\mathbf{E}^{(1)}))_i + \frac{\omega_1^2}{c^2}\epsilon_r E_i^{(1)} = -\frac{(\omega_2+\Omega)^2}{c^2}\chi_{ijkl}E_j^{(2)}u_{k,l} \tag{8}$$

Remarkably, electrostriction induces an optical volume force that drives the elastic wave equation according to[9]

$$\rho\frac{\partial^2 u_i}{\partial t^2} - (c_{ijkl}u_{k,l})_{,j} = T_{ij,j}^{\text{es}} \tag{9}$$

with the electrostriction stress tensor given by

$$T_{ij}^{\text{es}} = -\epsilon_0\chi_{klij}E_k^{(1)}E_l^{(2)*} \tag{10}$$

It is important to remark that the same photo-elastic tensor enters the non-linear polarization of Eq. (3) and the electrostriction stress tensor of Eq. (9). Actually, it is possible to consider an interaction Hamiltonian for the three-wave interaction so that the dynamical equations (7-9) can be derived from a single energy balance identity.[9–12] The possibility of energy transfer between a pump wave and signal optical wave, and an acoustic wave, is especially interesting in case of guided wave propagation, such as (stimulated) Brillouin light scattering in optical fibers.[8,13–15] The same formalism applies equally to other guided wave situations, such as photonic wires[16] or phoxonic crystal waveguides. It is also important to remark that the optical force is the divergence of the electrostriction stress tensor, i.e., $(F_{ES})_i = T_{ij,j}^{\text{es}}$. Because of the involved spatial derivative, it is expected that the optical force will be stronger for tightly confined optical fields, but also in case $|\mathbf{k}_1 - \mathbf{k}_2|$ is large; this is exactly the situation found with stimulated Brillouin scattering.

2.3 Optomechanical coupling and radiation pressure

Optomechanical coupling of light with the mechanical motion of nanoresonators has attracted a lot of attention recently,[17] accompanied with impressive demonstrations of optical cooling of laser mirrors[18] and nanomechanical resonators,[19,20] cavity optomechanics,[21] optomechanical coupling of optical waveguides and resonators, and much more. An important concept in this context is the optical gradient force that is exerted by the evanescent field of an optical wave (e.g., from a waveguide) on an optical resonator. In the case of coupled waveguides, Povinelli *et al.*[22,23] have proposed that the optical gradient force be given by

$$F = -\frac{1}{\omega}g_{OM}U \tag{11}$$

where U is the total system energy, ω is the eigenmode frequency of the resonant system, and

$$g_{OM} = \left.\frac{\partial\omega}{\partial x}\right|_{\mathbf{k}} \tag{12}$$

is the optomechanical coupling. This last quantity is the variation of the resonant frequency when the "relevant" distance in the system varies; for coupled waveguides, this distance is the length of the air gap separating them, while spatial phase matching is achieved. The derivation of Eq. (11) involves the idea that a change in internal energy of the coupled system should equal the work done by the mechanical force. This *ad hoc* approach has been practically successful, but gives no clue as to the relation between the optical field generating the force and the detailed vibrations induced inside the resonator (i.e., in contrast with an expression such as (9-10) that dictates the precise elastic wave dynamics). In the context of the optical excitation of mechanical motion of nanomechanical resonators, the optomechanical coupling defined by Eq. (11) is also used to quantify the coupling rate between the optical and the mechanical modes of the resonator.[20] In this picture, its value is an intrinsic characteristic of the resonator, independently of the external optical waveguide that couples photons in and out of the resonator. Since it depends on the field distributions of the optical and mechanical modes, it must be obtained via integration over a volume fully containing the resonator.

The classical way to compute the time-averaged force exerted on a dielectric rigid medium surrounded by a vacuum is from the Maxwell (or rather Minkowski) stress tensor, written here for an isotropic dielectric,[24]

$$T_{ij}^{ms} = \epsilon(E_i E_j - \delta_{ij} E^2/2) + (B_i B_j - \delta_{ij} B^2/2)/\mu \tag{13}$$

With σ a closed surface enclosing volume V containing the dielectric medium, the total force is

$$\mathbf{F} = \int_V dV \,\mathrm{div}(T^{ms}) = \int_\sigma T^{ms} \cdot \mathbf{n} d\sigma \tag{14}$$

The statement (14) is for the total force acting on the resonator. It is tempting to try and give a physical significance to the quantities appearing under the integral signs. At first sight, we could say that T^{ms} is a local stress tensor that could be used as a right-hand side for the elastic dynamics equation (9), for instance. This interpretation, however, has been argued against by many authors since one century, for many different reasons connected either with the validity of the mathematical derivation, or the deep difference between microscopic and macroscopic electromagnetic fields. As a simple rationale, we will object that if that interpretation were valid, then it should describe electrostriction inside materials, which is contrary to experience. A second possibility is to give a physical meaning to $T^{ms} \cdot \mathbf{n}$ as the surface force (or pressure) acting on the resonator boundaries, while volume stress is given by electrostriction.[16,25] Whereas radiation pressure on a rigid solid is firmly grounded, its action on an elastic solid is not obvious to this author, or at least he has the theoretical provision that in classical derivations of the Maxwell stress tensor,[24] the deformations of the elastic medium are not considered. The direct approaches underlined in the next section escape this criticism.

2.4 Moving boundaries

The diffraction of optical waves by surface acoustic waves has attracted some attention in the past,[26,27] though it remains dominated by bulk acousto-optics in most practical systems. The situation is similar for Brillouin (or more generally inelastic) light scattering at the surface of materials as compared to within materials. It has thus long been observed that the photo-elastic effect could be complemented by the normal displacement of the surface forming a shallow surface diffraction grating. If we now consider a small resonator or a waveguide subject to vibrations, we can estimate its acousto-optic properties via Eq. (5), by considering explicitly that the dielectric constant changes slowly in time.

$$\nabla \times (\nabla \times \mathbf{E}) + \frac{1}{c^2} \frac{\partial^2 (\epsilon_r(\mathbf{r}, t)\mathbf{E})}{\partial t^2} = -\mu_0 \frac{\partial^2 \mathbf{P}}{\partial t^2} \tag{15}$$

There is a factor of 10^5 between elastic frequencies (up to the GHz range) as compared to optical frequencies (a few 100 THz), so that at any particular instant in time τ the optical wave can be described as seeing a deformed but static medium, described by the spatial distribution of $\epsilon_r(\mathbf{r}, \tau)$. By decomposing the period T of the elastic wave in a number of snapshots, the induced effect on optical propagation can be straightforwardly computed.[28] If a resonator is considered, then a quantity similar to the optomechanical coupling will be obtained, but with the relevant distance being the measure of the deformation of the resonator. If a waveguide is considered, then the quantity $\Delta n/n \approx \lambda/k \times \partial k/\partial |u|$ can be monitored instead. It is apparent that the induced modulation

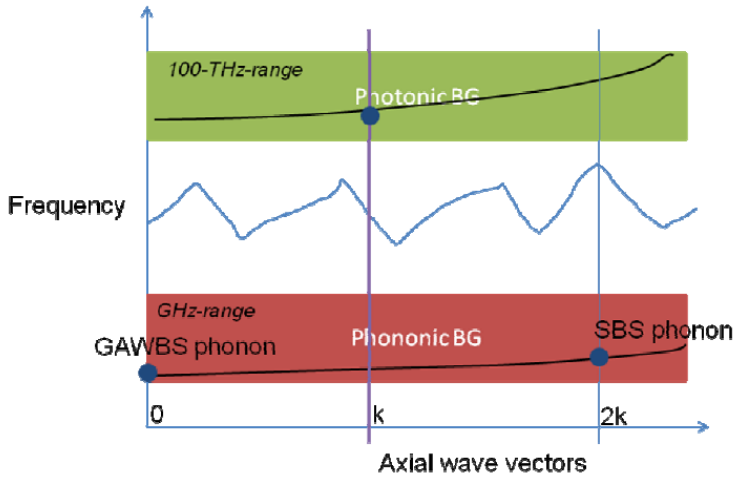

Figure 1. Sketch of a simultaneous photonic (green color) and phononic (red color) band gap in a phoxonic crystal. By adding a defect-line along the crystal, a phoxonic waveguide is formed, the dispersion of which is represented by the solid black line within each band gap. Considering the interaction of guided optical and elastic waves along the defect, two interaction cases are indicated: GAWBS and SBS (see text for definitions).

scales with the displacement of the boundaries, whereas the photo-elastic effect scales with the strain inside the medium. This suggests a cross-over between the two effects depending on the dimensions of the cross-section of the waveguide, as was found by Rakich *et al.* for nanoscale rectangular waveguides.[16]

3. A TENTATIVE PHOXONIC CRYSTAL ZOOLOGY

3.1 Photon-phonon interaction in a phoxonic band gap

As we outlined in the introduction, by phoxonic crystal we mean a periodic nanostructure that presents simultaneously a photonic and a phononic band gap. We have already stressed the huge difference in optical and elastic frequencies, about a factor of 10^5 in favor of the former. There is, however, no reason not to consider commensurate wavelengths for them, with the consequence that the spatial variations of optical and elastic wave can be engineered to be matched in phoxonic crystals. This situation is sketched in Figure 1.

It is well known that adding defects in an otherwise perfectly periodic structure allows one to form cavities and waveguides. In the case of cavities, the fields can be tightly confined by the surrounding periodic structure, leading in principle to the possibility of maximizing the optical and elastic modal overlap such as required by optomechanical coupling. In the dispersion diagram of Fig. 1, the cavity would have an optical resonant frequency appearing inside the photonic band gap and an elastic resonant frequency appearing inside the phononic band gap; both would appear as flat horizontal lines. In the case of waveguides, confinement is achieved in the cross-section while the waves are left free to propagate along the waveguide axis. While this would appear are less confining than a cavity, a waveguide authorizes to benefit – or not – from spatial phase-matching over long interaction lengths. Two cases can be considered for the interaction of 2 photons and one phonon, as in Brillouin scattering. In the forward interaction, the two photons are co-propagating; because $\Omega = \omega_1 - \omega_2 << \omega_1$, $K_z = k_{1z} - k_{2z} \approx 0$. In optical fibers, this case is known as guided acoustic-wave Brillouin scattering (GAWBS).[29–32] Because $K_z \approx 0$, the elastic modes involved are resonances of the cross-section of the waveguide. In the backward interaction, the two photons are contra-propagative, resulting in $K_z = k_{1z} + |k_{2z}| \approx 2k_{1z}$ (the elastic wavelength is twice the optical wavelength in the material). Stimulated Brillouin scattering (SBS) can then result, with both the elastic wave and the diffracted optical wave growing exponentially as they gain energy from the optical pump. The efficiency of the transfer for a given propagation length is also inversely proportional to the optical group velocity, making the occurrence of low group velocities particularly interesting.[33] Of course, in practice care has to be taken that propagation losses are not also strongly increased for low group velocities.

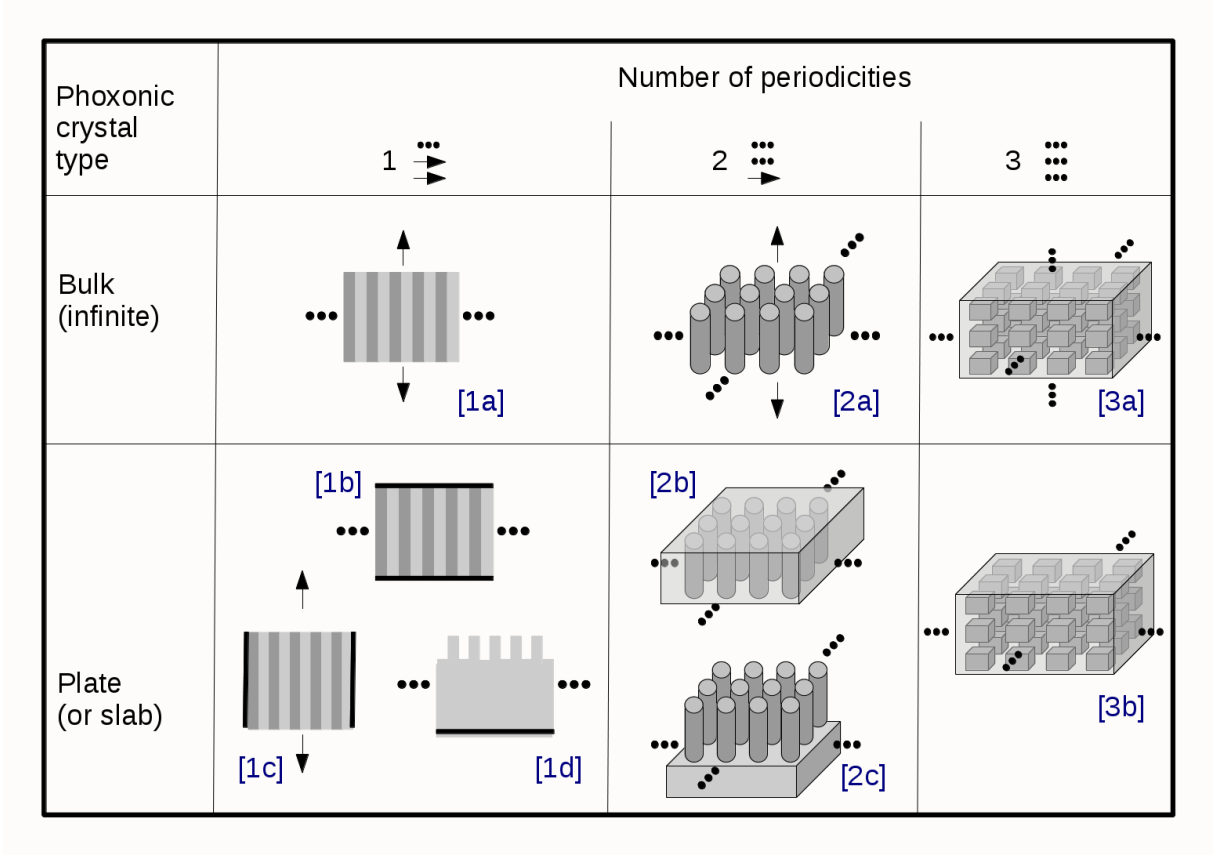

Figure 2. A tentative classification of possible phoxonic crystal types is shown. The classification is not exhaustive. Two basic geometries are considered: infinite crystals supporting bulk waves only and slab or membrane geometries that are amenable to technological fabrication with periodicities of a few 100 nm. Arrows (dots) represent infinite directions without (with) periodicity.

3.2 Types of phoxonic crystals

Figure 2 attempts a classification of the different types of geometries that have been considered – or not yet – for achieving phoxonic band gaps. The classification mainly distinguishes between the number of periodicities and whether hard boundary conditions terminating the crystal have to be taken into account, as with slabs or membranes.

The first and simplest type, [1a], is infinite and 1-periodic, and could be named a super-lattice. Dual photonic and phononic cavities were considered by Trigo et al.[34,35] before the concept of the phoxonic crystal was proposed. In these experiments, a phononic cavity (super-lattice with periodicity of a few nanometers) is enclosed within two photonic Bragg mirrors forming a Fabry-Perot cavity (a few 100 nm periodicity). Thz phonon generation was achieved in this structure from the optical pump and attributed to the strong confinement of both fields. Because phononic and photonic periodicities are different, the structure is not strictly speaking a phoxonic crystal as we defined before. Later, Papanikolaou et al.[28] proposed a theoretical study of a phoxonic crystal cavity formed of alternating silicon/silica layers. They showed that the strong simultaneous confinement could lead to nonlinear processes, and multi-phonon scattering. Photon-phonon coupling was a result of both the photo-elastic effect and the longitudinal vibration of the multilayer interfaces.

Type [2a] is the class of 2D infinite phoxonic crystals. It was for this geometry that maldovan *et al.* introduced the phoxonic crystal concept.[1,36] Specifically, they considered the case of air holes in silicon and the converse situation, silicon pillars in air. They found the first case is the most promising, because of the phononic properties of the structure. Sadat-Saleh *et al.*[2] performed an comprehensive search of phoxonic band gaps for a 2D infinite

crystal of air holes in lithium niobate. Lithium niobate has refractive indices smaller than silicon (around 2.2 instead of about 3.6), which renders the existence of complete photonic band gaps more difficult; the phononic crystal properties, however, are not strongly affected by the change in material, because the elastic contrast is given by the free boundaries of the holes, while optical fields extend in the air filling the holes as well as in the matrix. 2D infinite phoxonic crystals have remained idealized structures not directly accessible to experiments. There is a very close solution, however, that is provided by photonic crystal fibers (PCF).[37] It was shown in particular that PCF can support simultaneously phononic band gaps, and a nanostructure was proposed to act against SBS in silica optical PCF.[14] Trapped phonons within the solid core of a PCF were predicted and observed experimentally.[37] The phononic properties of guided phonons in PCF have been explored,[14,38,39] especially with respect to Brillouin effects.[31,32,40–42]

The class of 3D infinite phoxonic crystals (type [3a]) has not been considered in detail yet. Papanikolaou *et al.*[43] have predicted complete phoxonic band gaps in metallodielectric phoxonic crystals. Akimov *et al.* have performed pump-probe experiments on a colloidal 3D photonic crystal; they have observed efficient modulation of light with frequency close to a photonic band edge.[44] A demonstration of a true 3D phoxonic band gap, however, still has to be performed.

The case of 2D slab phoxonic crystals has received quite a lot of attention recently, for the reason that such structures are achievable by microelectronic-type nanotechnologies, e.g., silicon technology. Arrays of holes in a membrane (type [2b]) were investigated independently by Mohammadi *et al.* and Pennec *at al.*[3,45] It is observed that phoxonic crystal slab structures are achievable with silicon, but that the ratio of hole diameter to pitch has to be rather large. Square-lattice and honeycomb-lattice (or graphene) phoxonic crystals provide suitable phoxonic band gaps in theory, but not the hexagonal-lattice that is classical with photonic crystal slabs. This last property is imposed by the conditions for occurrence of phononic band gaps. The slab thickness, also, is a critical parameter, with its optimal value being roughly half the pitch. Safavi-Naeimi et al. have obtained quite different optimal designs using mass-and-spring type structures instead of holes.[46] Another alternative phoxonic crystal slab structure is the pillars on membrane structure[47] (type [2c]). Finally, a promising phoxonic crystal structure not depicted in Fig. 2 is the 1D phoxonic crystal strip.[48]

4. CONCLUSION

In this paper, we have attempted a description of phoxonic crystals in relation with their application to photon-phonon interaction. It is clear that quite a number of different nanostructures presenting simultaneous photonic and phononic band gaps are possible, yet they remain difficult to design and optimize. The main interests of such structures are (i) the possibility to confine simultaneously optical and elastic waves in cavities and waveguides, and (ii) the possibility to engineer the photonic and phononic dispersion relations of waveguides. A variety of coupling mechanisms are available for exaltation or quenching, including classical photo-elastic coupling as a volume interaction effect, and couplings introduced by moving boundaries or resonator dimensions. The physics and the applications of phoxonic crystals have not yet been explored fully, nor the very effect demonstrated experimentally with full evidence. As a consequence, there are still many opportunities for researchers in the field.

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