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Nanomechanical control of optical field and quality factor in photonic crystal structures

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Actively controlling the properties of localized optical modes is crucial for cavity quantum electrodynamics experiments. While several methods to tune the optical frequency have been demonstrated, the possibility of controlling the shape of the modes has scarcely been investigated. Yet an active manipulation of the mode pattern would allow direct control of the mode volume and the quality factor and therefore of the radiative processes. In this work, we propose and demonstrate a nano-optoelectromechanical device in which a mechanical displacement affects the spatial pattern of the electromagnetic field. The device is based on a double-membrane photonic crystal waveguide which, upon bending, creates a spatial modulation of the effective refractive index, resulting in an effective potential well or antiwell for the optical modes. The change in the field pattern drastically affects the optical losses: large modulations of the quality factors and dissipative coupling rates larger than 1 GHz/nm are predicted by calculations and confirmed by experiments. This concept opens new avenues in solid-state cavity quantum electrodynamics in which the field, instead of the frequency, is coupled to the mechanical motion.

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I. INTRODUCTION

The active manipulation of localized optical modes lies at the heart of many groundbreaking experiments in cavity quantum electrodynamics and studies of light-matter interaction, in particular because it allows on-demand control of the atom-photon interaction. So far, most of the efforts have been devoted to the tuning of the optical frequency in order to change the spectral overlap between the mode and a quantum emitter [1,2]. Several tuning methods, mainly based on changing the refractive index of the material surrounding the cavity [3–8], have been proposed. Additionally, frequency tuning can be obtained by a mechanical perturbation of the cavity structure itself, realized by, e.g., approaching and perturbing the optical near field by an external object, such as a tip [9,10], or changing the distance between different parts of the cavity. In this regard, micro- and nano-optoelectromechanical systems (NOEMS) have often been used as a reliable solution to implement spectral control on devices such as vertical-cavity surface-emitting lasers [11] and photonic crystal (PhC) cavities [12–18].

In all the approaches mentioned above the perturbation leads to a change in the optical frequency, while the spatial distribution of the field is only weakly altered. The opposite effect, i.e., the possibility of altering the spatial pattern of a cavity mode, has been recently considered in a few studies [19–22], including the demonstration of on-demand PhC cavity creation [23], and is quite attractive in various contexts. Indeed, the amplitude of the mode field at an emitter’s position directly determines the emitter-field coupling rate and affects radiative processes such as spontaneous emission, Rabi oscillations, and optical gain. Moreover, modifications of the field profile can lead to a change in the optical mode volume and the quality $Q$ factor, which also determine the strength of the light-matter interaction. If the perturbation is implemented mechanically, the resulting mechanical control of the optical losses can have important applications in the context of $Q$-switched lasers [24] and dissipative optomechanics [25]. In the latter, the change in the optical loss rate $\kappa$ is usually quantified through the dissipative coupling rate $g_\kappa \equiv \partial \kappa / \partial x$, where $x$ parametrizes the mechanical displacement. Previous approaches for the mechanical control of the optical loss have mainly relied on altering the external losses by controlling the distance between the cavity and an external reservoir [26–29] and/or altering the internal losses by displacing small parts of the cavity [27,29–33].

In electromechanical implementations, modest variations of the $Q$ factors ($\Delta Q / Q < 1$), corresponding to $|g_\kappa|/2\pi = 20–200$ MHz/nm [34], have been obtained so far [26,31].

In this work, we propose and experimentally demonstrate a concept to control the spatial pattern of a localized optical mode by a mechanical perturbation. The device is composed of two coupled PhC waveguides (PhCWGs) embedded in GaAs slabs [Fig. 1(a)] and separated by a small air gap. A bending of the membranes [Fig. 1(d)] creates a local modulation of the effective refractive index, which in turn controls the mode localization. In particular, an effective potential well or antiwell can be realized in our system, according to the symmetries of the modes considered. As a consequence, the optical fields get strongly localized either at the center or at the edges of the structure. The changes in the field profiles are very pronounced: for the case of the antiwell potential a mechanical displacement smaller than 10 nm is sufficient to completely reshape the pattern of the fundamental mode, so that the center of the waveguide evolves from being a point of maximum to a point of minimum of the field. Importantly, varying the field profile

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FIG. 1. Field and quality factor control in double-membrane NOEMS. (a) Schematic of the double-membrane structure, with a sketch of the symmetric (blue line) and antisymmetric (red line) field profiles. (b) Dispersion curves of the fundamental guided modes of a single-membrane PhCWG (green line) and of a double-membrane PhCWG for the symmetric (blue line) and antisymmetric (red line) modes. (c) In-plane field profiles (y component of the electric field) of the first three localized modes in a single-membrane PhCWG. (d) Sketch of the working principle of the device: when the membranes are bent towards each other, the S modes (only the first one is shown here) localize at the center of the waveguide. (e) and (f) SEM micrographs of a fabricated device. (g) Calculated downward displacement of the top membrane of the device shown in (e) under a voltage of 4 V applied between the membranes. The green dashed line indicates the position and the orientation of the PhC waveguide. (h) Distance between the two membranes along the waveguide direction for the maximum bending achievable in the real device [green curve, cut from (g) along the green dashed line] and for the case in which the edges are clamped (blue line), as shown in (d).

The device is schematically depicted in Fig. 1(a). It is composed of two parallel GaAs membranes containing two identical PhCWGs. The membrane thicknesses are 170 nm, the gap between the membranes is $z_0 = 200$ nm, the lattice constant of the PhC is $a = 370$ nm, and the filling factor of the holes is 0.35. The hybridization of the fundamental guided modes of the two isolated PhCWGs leads to the formation of two coupled guided modes, with symmetric (S) and antisymmetric (AS) vertical distributions of the y component of the electric field [Fig. 1(a), blue and red lines, respectively]. The dispersion curves of the S and AS guided modes [blue and red lines, respectively, in Fig. 1(b)] are spectrally shifted with respect to the dispersion curve of the fundamental guided mode of the single PhCWG [Fig. 1(b), green line]. The propagation of guided modes in coupled waveguides can be described in terms of effective refractive indices (ERIs) $\tilde{n}_i$ (i = S, AS) that depend on the materials, the membrane thicknesses, and the intermembrane gap [35,36]. In our structure, when the gap decreases, $\tilde{n}_S$ increases, while $\tilde{n}_{AS}$ decreases (see Supplemental Material (SM) [37]). We now consider the double-membrane PhCWG to be abruptly terminated: if the membranes are parallel [Fig. 1(a)], $\tilde{n}_i$ is uniform along the waveguide, and the structure can be described as a Fabry-Pérot cavity with weakly reflecting mirrors at the waveguide ends, caused by the large group-index mismatch. This gives rise, for both the S and AS guided modes, to a discrete set of localized modes in the slow-light region [i.e., $ka/2\pi \approx 0.5$; see fig. 1(b)]. The in-plane electric fields of the localized modes stem from the Bloch mode of the infinite waveguide, modulated by slowly varying envelopes [Fig. 1(c)]. We now consider the situation in which the membranes are bent in such a way that the intermembrane gap changes along the waveguide direction [Fig. 1(d)]. If the maximum bending is much smaller than the membrane’s length, we can describe the structure as a waveguide with a gentle spatial modulation of the ERIs. In particular, for a bending that is small with respect to the initial distance, the ERI depends exponentially on the local value of the air gap [38] and can be described by the formula

$$\tilde{n}_i(x, \Delta z) = \tilde{n}_\infty + \left(\tilde{n}_{i,0} - \tilde{n}_\infty\right)\exp[\gamma \Delta z \xi(x)],$$

where $\gamma$ is a constant and $\xi(x)$ is the spatial modulation of the gap.
where $\tilde{n}_\infty$ and $\tilde{n}_{i,2}$ ($i = \text{S, AS}$) are the ERIs of the single-slab waveguide and those of a double-slab waveguide with constant gap $\zeta_0$, respectively, $\gamma \equiv (\omega/c)/\sqrt{\tilde{n}_\infty^2 - 1}$ is the spatial decay rate of the single-slab electric field along the $z$ direction, $\omega$ is the angular optical frequency, and $c$ is the speed of light in vacuum. $\Delta z$ denotes the maximum variation of the gap [see Fig. 1(d)], and the function $\xi(\omega)$, normalized such that $\max(\xi) = 1$, describes the bending of the membranes along the $x$ direction [37]. The distance between the two slabs in each point is then given by $\zeta_0 - \Delta z \xi(\omega)$. For the structure considered here, the ERI of the S modes is larger than that of the single-slab mode ($\tilde{n}_{\text{S},2} > \tilde{n}_\infty$), as shown in Fig. S4 in the SM [37]. Therefore, when the deformation function $\xi(\omega)$ has a local maximum (i.e., where there are two slabs closer), the ERI of the S modes becomes higher in that region compared to its surroundings, promoting the localization of light. Conversely, $\tilde{n}_{\text{AS},2} < \tilde{n}_\infty$, and therefore, the ERI of the AS modes is decreased in the local maxima of $\xi(\omega)$, resulting in an antiwell which “pushes away” the AS modes. Thus, for a quasiperiodic bending profile, like the one shown in Fig. 1(d) (obtainable by applying a uniform force between the two clamped membranes), a strong localization of the S modes at the center of the waveguide is expected, while the AS modes will be pushed towards the waveguide’s edges. Moreover, as the optical losses of these modes are mainly due to the in-plane leakages at the waveguide-bulk interfaces, such variations of the field intensity at the waveguide edges is expected to strongly increase (decrease) the $Q$ factor of the S (AS) modes.

More useful insights into the behavior of the optical modes of this device can be obtained within the effective refractive index approximation [35,36] and by utilizing a one-dimensional envelope-function model [39,40]. The device is here described as a two-dimensional PhCWG made of air holes embedded in a material with a refractive index given by Eq. (1). The effect of the mechanical deformation can therefore be described by a perturbation of the periodic dielectric function $\epsilon(\omega)$ such that $\tilde{\epsilon}(\omega) = \Delta \epsilon(\omega) \epsilon(\omega)$ is the new dielectric function and $\Delta \epsilon(\omega) \equiv \tilde{n}_{i,2}^2(\omega, \Delta z)/\tilde{n}_{i,2}^2(\omega) (i = \text{S, AS})$, where the explicit dependence of $\Delta \epsilon$ on $\Delta z$ has been omitted for brevity. Moreover, in the following we omit the subscript $\omega$ from all the quantities since all formulas apply separately to the S and AS modes. It can be shown [39,40] that, for a generic slowly varying perturbation $\Delta \epsilon(\omega)$, the electric field of an unperturbed band-edge guided mode $\mathbf{E}_0(x, \omega)$ is modified into $\mathbf{E}_0(x, \omega, \Delta \epsilon(\omega)) \approx \epsilon^{k_{\text{B}} 2} \mathbf{u}_{k_{\text{B}}}(x, \omega)$, where $\mathbf{u}_{k_{\text{B}}}(x, \omega)$ is the periodic part of the Bloch function at $k_{\text{B}} = \pi/a$, $\alpha$ is the frequency of the perturbed mode, and we have introduced an adimensional function $\Gamma_\alpha(x)$. This latter function describes the modulations of the electric field envelope due to the perturbation, and it satisfies the equation [39,40]

$$
- \frac{c^2}{2m^*} \frac{d^2}{dx^2} + \omega^2(k_{\text{B}}) \Gamma_\alpha(x) = \alpha^2 \Delta \epsilon(\omega) \Gamma_\alpha(x),
$$

(2)

where $c$ is the speed of light, $\omega(k_{\text{B}})$ is the band-edge frequency of the unperturbed guided mode, and we introduced the adimensional effective mass $(m^*)^{-1} \equiv \frac{1}{\epsilon^2} \frac{d^2\epsilon(k_{\text{B}})}{dk_{\text{B}}^2}$. The derivation of Eq. (2) and the approximations involved are discussed in the SM [37]. Equation (2) is a generalized eigenvalue equation whose set of solutions provides the frequencies $\alpha$ and the field profiles $\Gamma_\alpha(x)$ of the modes created by the perturbation $\Delta \epsilon(\omega)$. As shown in the SM [37], for our device the function $\Delta \epsilon(\omega)$ has quadratic behavior near the center of the waveguide ($\omega \approx 0$). We can therefore expand $\Delta \epsilon(\omega) \approx \Delta \epsilon_0 + \Delta \epsilon_2 \omega^2$. By substituting this expression in Eq. (2) we obtain

$$
[ - \frac{c^2}{2m^*} \frac{d^2}{dx^2} + \omega^2(k_{\text{B}}) ] \Gamma_\alpha(x) = \alpha^2 \Delta \epsilon_0 \Gamma_\alpha(x) + \alpha^2 \Delta \epsilon_2 \omega^2 \Gamma_\alpha(x),
$$

(3)

which, depending on the sign of $\Delta \epsilon_2$, is mathematically equivalent to the Schrödinger equation of a generalized [i.e., the potential energy $V(x) = -\alpha^2 \Delta \epsilon_2 x^2$ depends on the “eigenenergy” $\alpha^2$] quantum harmonic or reversed oscillator. For the S modes, $\Delta \epsilon_2$ is negative, and it decreases with increasing membrane bending (see Fig. S4(b) in the SM [37]). Equation (3) describes therefore a generalized quantum harmonic oscillator [41,42], with an attractive potential energy at the center of the waveguide that increases with bending. For the AS modes the term $\Delta \epsilon_2$ is positive, and it increases with increasing membrane bending. Equation (3) describes in this case a generalized quantum reversed oscillator, in which the modes feel a repulsive potential at the origin $x = 0$ and are therefore pushed towards larger values of $|x|$.

**B. Numerical analysis**

While the analytical model discussed above provides a semiqualitative description of the control of the optical fields in our device, its accuracy in predicting the field and loss modulations is hindered by the approximations involved, as discussed in detail in the SM [37]. We have therefore utilized full three-dimensional finite-element-method (FEM) simulations [43] to verify the magnitudes of the predicted effects. Figures 2(a) and 2(b) show the calculated (FEM) in-plane profiles of the electric field ($y$ component) of the first two S modes for a PhCWG length of $L_{\text{WG}} = 28a \approx 10.5 \mu m$ and a quasiperiodic bending of the membranes [Fig. 1(d)]. The bending profile is calculated analytically by solving the Euler-Bernoulli beam equation for a doubly clamped beam under uniform applied force [37]. The bending is parametrized by the decrease $\Delta \zeta$ of the gap in the center of the waveguide. As expected from the previous discussion, the S modes get strongly localized at the waveguide’s center for increasing bending [Figs. 2(a) and 2(b), from left to right], as a result of the local increase of the ERI (the field patterns for different displacements and for the other modes are shown in the SM [37]). Figures 2(c) and 2(d) show, respectively, the frequencies and the loss rates of the first four S modes versus $\Delta \zeta$. The frequencies of all modes decrease [Fig. 2(c)] due to the overall increase of the ERI. Moreover, the frequency spacings, initially strongly uneven, become uniform for large bending, as expected for the eigenmodes of a potential well, which can be approximated locally by a harmonic oscillator. Importantly, the loss rates of all the S modes [Fig. 2(d)] decrease very strongly (by more than two orders of magnitude for the first two modes) as a consequence of the reduced leakages at the waveguide edges. For $\Delta \zeta > 30 \text{ nm}$ the losses of the first two modes saturate and start to increase. We attribute this behavior
FIG. 2. Numerical analysis of the field localization in bent double-membrane waveguides. \( \Delta z \) denotes the variation of the distance between the centers of the two membranes. (a) and (b) In-plane field profiles (\( y \) component) of the first two S modes for no bending (parallel membranes, \( \Delta z = 0 \) nm) and a moderate bending (\( \Delta z = 30 \) nm), as indicated on top of each panel. The color scales of all plots are the same as in Fig. 1(c). (c) Frequencies of the first four S modes versus \( \Delta z \). (d) Optical loss rates of the first four S modes versus \( \Delta z \). The color legend is the same as in (c). The vertical axis on the right shows the corresponding values of the \( Q \) factor, calculated for an average normalized frequency of 0.277. (e)–(h) Same as in (a)–(d) for the AS modes. The \( Q \) factor scale in (h) is calculated for an average normalized frequency of 0.2897. (i) Mode volume \( V \) (solid lines) and \( Q/V \) ratio (dashed lines) of the S modes for increasing bending. (j) Modulus square of the Fourier transform of the in-plane profile of the electric field (\( y \) component) of the first S mode for \( \Delta z = 0 \) nm and \( \Delta z = 60 \) nm. The red circle denotes the light cone. (k) Modulus square of the Fourier transform of the in-plane electric field (sum of the \( x \) and \( y \) components) of the first S mode integrated over the light cone versus the displacement \( \Delta z \).

Remarkably, we note that a moderate bending can provide localized modes with extremely high optical \( Q \) factors (\( Q > 10^7 \)), comparable with the ones obtainable in PhC cavities where the position and/or radius of each hole is carefully optimized to maximize the \( Q \) [45–49]. Moreover, the field localization of the S modes leads to a twofold reduction of the mode volume \( V \) [Fig. 2(i), left axis], which further increases the ratio \( Q/V \), a typical figure of merit for cavity QED experiments and, in particular, for Purcell enhancement. As shown in Fig. 2(ii) (right axis), our structure can attain values of \( Q/V \) up to 10^7(\( \lambda^3/n^3 \))^{-1} that are comparable to those theoretically predicted in properly optimized PhC cavities [45,48] or hybrid plasmonic-PhC designs [50]. We notice that, while a bending of \( \Delta z \geq 30 \) nm is required to
obtain the lowest absolute values of the losses, the dissipative coupling rates are already large for small displacements around the mechanical equilibrium. The dissipative coupling rates of the first four S modes, calculated at \( \Delta z = 0 \), are \( \{ g_x^{(1)}, g_x^{(2)}, g_x^{(3)}, g_x^{(4)} \}/2\pi = \{-0.82, 3.19, 4.31, 3.71\} \) GHz/nm, which are among the largest dissipative coupling rates predicted so far for optical modes in the visible and infrared range.

In Figs. 2(e)–2(h) we show the same calculations as in Figs. 2(a)–2(d) but for the AS modes. When the structure is bent, the field profiles of the first two AS modes are pushed towards the edges of the waveguide, and except for the opposite phase of the two lobes, the field patterns become identical. Similar behavior occurs for the third and fourth AS modes, as shown in the SM [37]. Accordingly, the frequencies of the first two AS modes become quasidegenerate with increasing bending [see Fig. 2(g) and its inset, red and black lines], and the same happens for the third and fourth modes (green and blue lines). Such a degeneracy is a consequence of the antwell potential. Importantly, due to the localization of the fields at the waveguide edges, the losses of all the AS modes strongly increase for large deformations [solid lines in Fig. 2(h)]. Large dissipative coupling rates, \( \{ g_x^{(1)}, g_x^{(2)} \}/2\pi = \{0.76, 1.46\} \) GHz/nm, are therefore also predicted for the AS modes at \( \Delta z = 0 \).

C. Experimental results

The double-membrane GaAs PhCWG samples were fabricated through several steps of optical and electron-beam lithography, similar to the process used for parallel double-membrane PhC cavities [18] and described in detail in the SM [37]. A layer of high-density InAs quantum dots is grown in the middle of the top membrane and used as an internal source of excitons. The bending of the structure is achieved electromechanically by applying a voltage between the adjacent faces of the two membranes. The maximum decrease in the intermembrane gap obtainable with this actuation scheme, known as the pull-in limit, is \( \Delta z_{\text{max}} = 1/3 \) \( z_0 \approx 67 \) nm in our geometry [37]. In the fabricated device, only the top PhCWG is mechanically compliant, and it is embedded in a four-arm bridge, while the bottom PhCWG is embedded in a fixed membrane [Figs. 1(e) and 1(f)]. This design allows us to strongly reduce the postfabrication vertical buckling [51] by etching small rectangular trenches close to the points where the four arms are connected to the substrate, as suggested in a recent work [51]. However, compared to the ideal double-bridge design [Fig. 1(d)], this design features two main differences. First, only the top membrane will bend under an applied intermembrane force. Nevertheless, as long as the vertical displacement is small compared to the length of the bridge, we expect no fundamental difference with respect to the case in which both membranes move towards each other (as assumed in the calculations in Fig. 2). Second, the edges of the top waveguide are not clamped. Therefore, even if a quasiparabolic bending along the waveguide direction is obtained for large applied voltages [Fig. 1(g)], the difference between the bending of the center and that of the edges of the membrane is strongly reduced with respect to the ideal case in which the edges do not move [compare green and blue curves in Fig. 1(h)]. In particular, close to the pull-in condition, the bending difference is approximately 10 nm for the real structure, compared to 67 nm for the ideal one. As we discuss below and in the SM [37], this suboptimal bending profile represents one of the main practical limitations in the present implementation.

In Fig. 3(a) we show the photoluminescence spectra of the S modes in a device with lattice constant \( a = 385 \) nm for different applied voltages. Three peaks are visible in the spectrum when no electrostatic force is applied \([V_B = 0 \text{ V}, \text{dark blue shaded spectrum in Fig. 3(a)}]\). The highest-wavelength peak actually consists of two almost overlapping modes, which is more evident for larger applied voltages \( (V_B > 3 \text{ V}) \) when the two modes split apart. For each applied voltage we extracted the frequency and linewidth of the optical modes by performing a multiplet Lorentzian fit. The frequency spacing between the four modes is initially strongly uneven [see Fig. 3(b) at \( V_B = 0 \text{ V} \)]. As the voltage is progressively increased, the whole set of modes redshifts, and the mode spacings become almost uniform for large applied voltages, as expected for the energy levels of a harmoniclike potential. Moreover, as shown in Fig. 3(c), the optical loss rates of all four modes decrease for large applied voltages. For the first mode \([\kappa_1, \text{black squares in Fig. 3(c)}]\), a modulation of the loss rate of a factor of almost 2 is obtained, while for the second mode \([\kappa_2, \text{red circles in Fig. 3(c)}]\) the loss rate is modulated by a factor larger than 3 (calculated from the lowest to the highest point). A decrease in the loss rate of a factor of 2 is also observed for the third and fourth modes. The initial increase in the loss rates of the first two modes is likely due to the structural disorder and to the frequency degeneracy of the two modes, which, up to a voltage of 3 V, makes the estimation of the linewidths difficult. By comparing the measured spectral tuning of the first mode [Fig. 3(b), black line] with the simulated one, we can infer the vertical displacement of the center of the waveguide for each applied voltage, as indicated in the top axis of Figs. 3(b) and 3(c). This allows us to estimate the maximum dissipative coupling rates. In particular, by performing a fit of the data shown in Fig. 3(c) for large applied voltages (specifically, for \( V_B > 3 \text{ V} \) for the first mode and for \( V_B > 3.4 \text{ V} \) for other modes) we obtain the values \( \{ g_x^{(1)}, g_x^{(2)}, g_x^{(3)}, g_x^{(4)} \}/2\pi = \{2.4 \pm 0.7, 6.2 \pm 1.0, 2.0 \pm 0.1, 2.3 \pm 0.1\} \) GHz/nm, which are in the same range as those calculated for the ideal structure in Fig. 2.

In Figs. 3(d)–3(f) we show the results of the electrostatic tuning of the AS modes of a different device, which is nominally identical to the first one (apart for the lattice constant, which is \( a = 380 \) nm) but fabricated in a different run. Two modes are clearly visible and well separated in the spectrum when no electrostatic force is applied [bottom plot in Fig. 3(d)]. When the structure is bent [Fig. 3(d)], the two modes blueshift, and their spectral spacing quickly decreases until, at about 2.6 V, the modes completely overlap. Moreover, the optical losses of both modes strongly increase, resulting in a loss modulation of more than a factor of 3 for the first mode and more than a factor of 5 for the second mode [Fig. 3(f)]. For voltages higher than 2.95 V, the mode intensity becomes too low, and the linewidth cannot be estimated accurately. Both the spectral merging of the two modes and the loss increase match the expected behavior shown in Fig. 2(g) (red and black lines). Like for the S modes, we estimate the membrane displacement [top axis of Figs. 3(e) and 3(f)] by comparing the experimental frequency tuning of the first mode with the
FIG. 3. Control of the optical losses in photonic crystal waveguides. Positive voltages correspond to attractive electrostatic force. (a) Spectrum of the S modes for increasing applied voltages. The plots are vertically shifted for clarity. (b) Frequencies and (c) optical loss rates of the four S modes versus the applied voltage, extracted from (a). The top horizontal axis shows the calibrated displacement of the membrane’s center, obtained by comparing the measured frequency tuning of the first mode with the calculated data. (d)–(f) Same as in (a)–(c), but for the AS modes of a different device.

simulated one. From this calibrated displacement, we calculate the maximum dissipative coupling rates of the two AS modes, \( \{ g^{(1)}_s, g^{(2)}_s \} / 2\pi = \{ 3.3 \pm 1.7, 7.9 \pm 3.1 \} \text{ GHz/nm} \).

III. DISCUSSION

While the observed trends match the theoretical results, the measured optical loss rates deviate considerably from the theoretical values, both in terms of their initial values (i.e., when the structure is not bent) and in terms of the total loss modulation, which is expected to be more than two orders of magnitude for the first two S and AS modes. A primary cause of this deviation is constituted by the structural disorder in the PhC pattern and by the residual surface roughness due to the fabrication. Both these imperfections introduce additional loss channels, which can become the dominant loss mechanism and therefore hinder the expected loss modulation. Moreover, in the fabricated devices the obtainable bending profile is far from the ideal one [compare blue and green lines in Fig. 1(h)] because the edges of the waveguides are also displaced downward. This reduces the contrast between the ERI of the center and that of the edges of the waveguide and therefore limits the field modulation. We quantitatively investigated the effect of the structural disorder and the nonoptimized bending separately by performing a set of numerical simulations [37]. The structural disorder can be modeled by introducing in the FEM simulations random fluctuations of the positions and radii of the holes. Our calculations (Fig. S14 in the SM) show that an amount of structural disorder comparable to that observed in the fabricated devices can limit the maximum \( Q \) factor achievable in a bent waveguide to values of \( 10^4 \)–\( 10^5 \), thus reducing the loss modulations calculated for a nondisordered PhC (in particular for the first two S/AS modes). To estimate the impact of the suboptimal bending profile, we performed FEM simulations in which the intermembrane spacing is varied according to the bending profile of the real device [such as the one in Fig. 1(g)], calculated for different applied voltages. The results (Fig. S13 in the SM) indicate that a realistic bending profile can substantially reduce the loss modulations, particularly for the S modes, for which the maximum \( Q \) factor enhancements shown in Fig. 2(d) are reduced by a factor of about 10. Finally, we performed simulations in which both the structural disorder and the suboptimal bending profile are taken into account (Fig. 4). The calculations show that, while a clear decrease in the optical losses is still obtained for large values of \( \Delta z \), the loss modulations are less than one order of magnitude [Fig. 4(b), averaged over four realizations of the disorder], in good agreement with the experimental data. We note that the dependence of the loss rate on the bending becomes nonmonotonic for some modes in the presence of imperfections, a feature also observed in the experiments [Fig. 3(c)]. The frequency spacings [Fig. 4(a)] become even for the first three modes for large displacements, while the spacing between the third and fourth modes remains larger, similar to what is observed in the experimental data [Fig. 3(b)].

The field patterns of the first two S modes (for a particular
realization) are quite affected by the structural disorder for \( \Delta z = 0 \) [Figs. 4(c) and 4(d), first row]. Nonetheless, for large mechanical deformations a reshaping and a localization of the field profiles are clearly visible. We emphasize that, despite these strong limitations in the current implementation of the device, the measured modulation of the optical losses is much larger than the one obtained in other electromechanical systems [26,31], demonstrating the potential of our method for the mechanical control of the \( Q \) factor in PhC structures.

Importantly, while the total modulation of the loss rate is reduced by the effects described above, we do not expect the dissipative coupling rate to be affected. Indeed, the variation of the loss rate is not altered by the presence of other loss channels. Moreover, large dissipative coupling rates are expected in the ideal device already for small displacements \( \Delta z \). Around the configuration in which the membranes are parallel to each other [see Figs. 2(d) and 2(h)]. Such moderate bending profiles can be reached in real devices for large applied voltages [e.g., Fig. 1(h), green line]. Indeed, the order of magnitude of the experimental values of the dissipative coupling rates at large voltages \((g_\kappa/2\pi = 1 - 10 \text{ GHz/nm})\) agrees well with the dissipative coupling rates calculated for the ideal device and for small displacement around \( \Delta z = 0 \).

Finally, we note that the displacement pattern induced by the applied voltage is rather similar to that of the fundamental flexural mechanical mode of the top membrane. Therefore, the values obtained for the dissipative coupling rates \( g_\kappa \) allow us to estimate the vacuum dissipative coupling rates (i.e., the loss rate modulation due to the presence of one phonon) achievable in our structure. The vacuum dissipative coupling rate is defined by \( g_{0,\kappa} \equiv g_\kappa x_{zpf} \), where \( x_{zpf} = \sqrt{\hbar/2M_\text{eff}\Omega} \) is the zero-point motion amplitude and \( M_\text{eff} \) and \( \Omega \) are the effective mass and the resonant frequency of the mechanical mode, respectively. By calculating numerically \( M_\text{eff} \) and \( \Omega \), we obtain a value of \( x_{zpf} = 9.6 \text{ fm} \), which leads to vacuum dissipative coupling rates in the range \( |g_{0,\kappa}|/2\pi = 10-100 \text{ kHz} \).

**IV. CONCLUSIONS**

We have proposed, theoretically discussed, and experimentally demonstrated a photonic-crystal-based NOEMS in which the mechanical bending of a membrane induces large variations of the spatial pattern of the optical fields. This is achieved by locally modulating the effective refractive index in a double-membrane photonic crystal waveguide, which results in the creation of a potential well or antwell for the optical field. As a consequence of moving the optical field far away or close to the waveguide’s edges, large modulations of the optical losses are predicted by our calculations. Experimentally, we were able to measure modulations of the optical losses up to a factor larger than 5, much larger than the ones obtained with different electromechanical approaches [26,31,32], and dissipative coupling rates in the 1–10 GHz/nm range. We have shown that the performance of our devices is currently limited by the structural disorder of the PhC and by the suboptimal bending profile. The latter is due to the particular four-arm bridge design that we used to eliminate the postfabrication vertical buckling. Thus, the bending profile can be improved by either further optimizing the geometrical parameters of the four-arm bridge design or by working with shorter doubly clamped bridge designs [Fig. 1(d)] that are not affected by vertical buckling (Fig. S11); since shorter bridges are also stiffer, the latter option will require fabricating \( p-n \) junctions with larger breakdown voltages or implementing different bending mechanisms.

Importantly, this work demonstrates how the field pattern in a nanophotonic cavity can be controlled mechanically and opens the way to new regimes of optoelectromechanical coupling in which the field, instead of the frequency, is coupled to the mechanical motion. We note that, while in the present structure the mechanical motion produces a variation of mode frequency, field distribution, and loss rate, it is possible to design structures in which only the field is modulated upon an antisymmetric perturbation [20–22]. The mechanism investigated in this work could lead to several applications such as novel platforms for controlling the spontaneous emission based on altering the field amplitude at the emitter position rather than introducing a frequency detuning [19], establishing an efficient emitter-phonon coupling [21], and the mechanical control of the far-field emission and lasing.

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[34] The quoted values of the dissipative coupling rates have been inferred by the authors from Fig. 3(b) of Ref. [31] and Fig. 3(b) of Ref. [26].


