Three dimensional trapping of light with light in semiconductor planar microcavities

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Abstract

When light is confined in all three directions and in dimensions of the order of the light wavelength, full discretization of the photon spectra and distinctive phenomena occur, the Purcell effect and the inhibition of emission of atoms being two paradigmatic examples. Diverse solid-state devices that confine light in all three dimensions have been developed and applied. Typically the confinement volume, operating wavelength, and quality factor of these resonators are set by construction, and small variations of these characteristics with external perturbations are targeted for applications including light modulation and control. Here we describe full 3D light trapping, that is set and tuned by laser excitation in an all-optical scheme. The proposed device is based on a planar distributed Bragg reflector GaAs semiconductor microcavity operated at room temperature. Lateral confinement is generated with an attractive effective photonic potential tailored by the increase in the cavity spacer refractive index, induced by carriers photoexcited by a focused laser. Strong three dimensional trapping of light is evidenced by the laser-induced changes on the spectral, spatial, and k-space distribution of the emission. The dynamics of the laser induced photonic potential is studied using modulated optical excitation, highlighting the central role of thermal effects at the origin of the observed phenomena.
Resonant cavities are pervasive to many different natural phenomena and man made devices. Standing waves confined between reflective walls build up, localizing the wave energy and enhancing their interaction with other physical degrees of freedom. They are ubiquitous in the domain of optics, for example as the basic feedback mechanism of lasers\textsuperscript{1,2}, or in the most sensitive displacement sensors, such as the Fabry Perot cavities used for the detection of gravitational waves at LIGO\textsuperscript{3}. When the confinement is strong and in all three dimensions, and the size of the optical cavities is of the order of the light wavelength, full discretization of the photon spectra occurs\textsuperscript{2,4–6}, leading to distinctive phenomena as for example the inhibition or enhancement of the emission of atoms\textsuperscript{7–9}, and the Bose-Einstein condensation of strongly coupled matter and light particles (cavity polaritons)\textsuperscript{6,10}. In the NIR-visible spectral range, such three-dimensional photon trapping typically requires the development of confining structures based on nano- and micro-fabrication methods. Diverse techniques have been developed and applied, including the use of local defects in periodic photonic bandgap free-standing membranes\textsuperscript{11}, the introduction of traps through the lateral patterning of planar distributed Bragg reflector (DBR) cavities\textsuperscript{2,6,12}, or their full 3D structuring into microdisks and pillar resonators based on finely tuned ion etching methods\textsuperscript{4,5,13}. These techniques are already mature and robust, leading to high quality devices that have impacted enormously in the landscape of novel optical, optoelectronic and optomechanical phenomena in the last couple of decades. One key targeted application has been that of all-optical modulation and control of light. Being solid state devices, however, their flexibility regarding their confinement characteristics and operating wavelength is limited, since it is typically defined by design and construction. In the most standard approach, highly confining resonators are used with the ultra-narrow cavity mode subtly tuned by small changes in refractive index\textsuperscript{14,15}. A different situation would be attained if light trapping could not only be tuned but also set in an all-optical scheme. Such tunable 3D trapping of light is what we describe here.

Our approach to this task is based on a planar DBR semiconductor microcavity for the light confinement in one direction ($z$), plus a photon potential defined by focused laser excitation for the in-plane additional trapping (see the scheme in Fig. 1a). When a semiconductor is optically excited above or close to an electronic transition, a change of the
Figure 1: a) Sketch of the system, showing the generation of electron-hole pairs in the GaAs spacer and the resulting temperature distribution. The e-h pairs are confined in the spacer due to the Al-rich neighboring layers; their non-radiative recombination acts as a localized heat source. b) Representation of a Gaussian optical well and the resulting confined modes, which appear beneath the original continuous spectrum. c) Photoluminescence spectra as a function of incident power, ranging from 1.8 mW up to 22.2 mW. d) Optical potential well depth as a function of laser power. The black circles are experimental results, while the red line is a linear fit to the low power ($\leq 30$ mW) part of the data, and the blue dashed line is a quadratic fit to the whole data.

refractive index occurs. This can be due to the modification of the electronic bands or the accessible electronic states induced by the presence of excited carriers$^{16,17}$, or by the lattice heating produced by the relaxation of these carriers$^{18}$. If the change is positive, the local value of the cavity mode energy will be decreased and then light trapping will be possible. A conceptual scheme of the so-obtained 3D effective attractive photon potential is shown in Fig. 1 b.
In the transparency region of a semiconductor, and for a fixed wavelength, the refractive index is expected to increase if the band gap is reduced. A gap reduction could be produced by an increase of the local temperature generated by photocarrier induced laser heating. However, more complex nonlinearities related to photoexcited carrier correlations could in principle also be involved\textsuperscript{16}. For example, polariton energy-barriers (in contrast with traps) have been demonstrated as a result of the blueshift of the polariton branch due to polariton-polariton interactions in microcavities at low temperatures and in the strong coupling regime\textsuperscript{19–21}. Trapping can be induced, for example, by means of potential barriers generated with ring-like laser illumination, as reported for cold excitons in indirect quantum wells\textsuperscript{22}. Optically generated local traps in planar microcavities with attractive weak potentials have been shown to stabilize polariton condensates with a mechanism based on spatially confined gain\textsuperscript{23}, similar to gain guiding in semiconductor laser structures\textsuperscript{24}. Also, positive changes of the refractive index have been reported to affect photon lasing in GaAs microcavities when evolving from the excitonic regime toward the electron-hole plasma regime with increasing photoexcitation power\textsuperscript{25}. In contrast with these low-temperature carrier-related correlation effects, we will demonstrate that strong three dimensional trapping of photons can be achieved in planar semiconductor microcavities at room temperature through photothermal phenomena.

The studied sample consists of a high Q-factor $\lambda/2$ bulk-GaAs planar cavity enclosed by two DBRs consisting of alternating Ga$_{0.9}$Al$_{0.1}$As/Ga$_{0.05}$Al$_{0.95}$As $\lambda/4$ layers, 28 pairs on the bottom, and 24 on top. The reported experiments correspond to room-temperature photoluminescence microscopy measurements. The optical setup, described in the Supplemental Material, allows for the acquisition of the integrated emitted spectra, or alternatively the spatial or k-dispersion profile of the emission\textsuperscript{26}. The sample was excited with a 1.63 eV cw Ti-Sapphire laser, and a 20x objective lens with a numerical aperture of 0.3 was used for both excitation and light collection. Figure 1c presents a few examples of the photoluminescence spectra obtained when exciting with different incident powers, and with the optical cavity mode tuned below but not far from the GaAs gap ($\varepsilon_{gap} \sim 1.42$ eV at room temperature).

The low power spectrum in Fig. 1c (1.8 mW) is representative of the planar microcavity (1D confinement) situation. The planar cavity quality factor in the transparency region, below the GaAs gap, is around $1 \times 10^4$. This corresponds, for the emission normal to the structure, to a narrow peak of $\sim 150 \mu$eV width centered in this case at $\sim 1.3755$ eV. The
confined cavity mode emission along $z$ defines the low energy flank of this curve. Its flat continuation towards higher energies is determined by the in-plane parabolic dispersion of the planar cavity optical mode, which is collected within the numerical aperture of the microscope objective\textsuperscript{27}. Major and quite striking changes in the spectra occur when the focused laser power is increased (note the vertical logarithmic scale). Some conspicuous new features arise, namely a red shift of the emission, and the appearance of clear peaks contrasting with the essentially featureless low-power spectrum. As we argue next, these features evidence the full optical confinement within the Gaussian photon potential well induced by laser excitation.

Full three dimensional confinement of light should be reflected in the induced changes of the spectra (as manifested by the energy lowering and spectral discretization in Fig. 1c), but also by the corresponding modified spatial and k-space (angular) characteristics of the photoluminescence emission\textsuperscript{28}. Figure 2 presents the spatial (a, c) and angular (b, d) distributions of the emitted light when a low (a, b) or high (c, d) power excitation is applied. Again, the low power panels (1.7 mW) are representative of the planar (1D confinement) case. As expected, the spatial distribution reproduces the shape of the excitation Gaussian profile. Its in-plane k-dispersion in turn reflects the typical parabolic behavior of light confined in one dimension, but free to propagate in the other two (in-plane) directions. The situation dramatically changes when increasing the focused laser power (18 mW in Fig. 2). New fully discretized modes appear at lower energies, with spatial and in-plane k-dispersion strongly reminiscent of laterally confined modes, as seen e.g. in micropillar cavities\textsuperscript{4,5,29} or in laterally microstructured planar resonators\textsuperscript{6,30}. The spatial and spectral envelope of the observed modes in Fig. 2c nicely reflects the Gaussian photon potential trap induced by the focused laser.

In Fig. 3 we show the lateral distribution of the emitted light for 18 mW (high power) excitation. The top left graph corresponds to the total integrated emission, reflecting the homogeneous spatial distribution of the excitation beam. The spectrum displayed in the top-right panel corresponds to this emission when the spot is imaged centered on the spectrometer entrance slit. The vertical dotted line in this latter panel shows the energy of the unperturbed cavity mode, above which modes are not confined. The discrete peaks (1)-(5) correspond to the observed fully 3D confined photonic states. The six remaining graphs in this figure are constant energy sections that were selected to leave only one mode for each
cut (labeled with the same numbers 1 to 5), and the non-confined spectral region. Circularly symmetric modes confined within different radial distances are clearly identified.

The spectral and spatial distributions of the observed modes have the required information to model the light induced photon potential. The incident laser generates, by mediation of the photo-excited carriers, an optical potential well that confines light in an analogous way as that observed in a micropillar structure\textsuperscript{4}. Based on the excitation profile and symmetry

Figure 2: Photoluminescence spectra spatial (a, c) and angular (b, d) resolution. The colors chosen for the spatial (orange) and the angular (blue) images are different to help on their rapid identification. The sample was excited with $\sim 1.7\text{mW}$ (a, b) and $\sim 18\text{mW}$ (c, d).
Figure 3: Spatial distribution of the photoluminescence spectra for an excitation of \( \sim 18 \text{ mW} \). The top leftmost graph corresponds to the total emitted light. The red curve is the spectral distribution of the emitted light coming from the central line (\( Y = 0 \mu\text{m} \)) of the spot. The rest are different energy cuts (with increasing energy from left to right and top to bottom, labeled from 1 to 5). The dashed orange arrows and circles show the calculated energies and effective diameters of the confined modes, respectively.

of the studied system, a Gaussian well of finite depth is expected to be generated:

\[
\varepsilon = \varepsilon_{00} - \Delta \varepsilon_0 e^{-\frac{r_{eff}^2}{2\sigma^2}},
\]

where \( \varepsilon_{00} \) is the energy of the fundamental optical mode in the absence of a well (data
Figure 4: Top panel: Left) photoluminescence spectra as a function of frequency for a 20% duty cycle square-wave modulation. From blue to red, the curves correspond to frequencies ranging from 1 MHz to 1 Hz. Right) photoluminescence spectra as a function of the duty cycle for a square-wave modulation of 400 kHz. From blue to red, the curves correspond to duty cycles ranging from 20% to 80%, with steps of 5%. Bottom panel: Optical well depth as a function of the modulation period for an excitation of 33 mW. Inset) Optical well depth as a function of the duty cycle time for a modulation of 400 kHz. The circles correspond to the experimental results, while the red dotted curves were obtained from a thermal diffusion model. The latter considers a periodic heat source in the GaAs spacer and a constant temperature plane at a distance of 7.5 μm from the spacer. The thermal diffusivity used for the fit was α_T = 10 μm²/μs.

obtained from the experiment), Δε₀ the effective depth of the well (fit parameter alongside σ) and D_{eff} the effective diameter of the well for an energy ε. The problem can be significantly simplified by making the approximation that each confined mode of energy ε, interacts with an infinitely deep circular well of diameter D_{eff}. This is equivalent to considering that the confined modes present a node on the wall of the Gaussian well. In a infinitely deep circular potential well, the confined modes can be well represented with Bessel functions, and their
energies can be calculated as:

\[ \varepsilon = \sqrt{\varepsilon_0^2 + \frac{4\hbar^2 c^2 X_{ml}^2}{n_{cav}^2 D^2}}, \]  

where \( D \) is the diameter of the confinement potential, \( X_{ml} \) the m-th root of the Bessel function of order l, \( c \) the speed of light in vacuum, \( n_{cav} \) the effective index of refraction of the sample and \( \varepsilon_0 \) the energy of the mode if no confinement existed. The theoretical results fit the peaks of the observed spectra very well (see for example the arrows with the spectra in Fig. 3, and additional information in the Supplemental Material). This agreement allows to conclude that the laser-induced potential shape is essentially independent of the laser power (its Gaussian shape and lateral width \( \sigma \) remain unaltered), being only the well depth strongly sensitive to the excitation power (as shown in Fig. 1d). The trapping energy attains a quite significant value of \( \sim 13 \text{ meV} \) for \( \sim 80 \text{ mW} \) excitation, under the focusing conditions used, a value equivalent to 100 times the FWHM on the unperturbed planar cavity mode. The insensitivity of the photon potential width \( \sigma \) with laser power can be expected from the invariant laser spot size. The supra linear behavior of the photon potential depth dependence with laser power, on the other hand, requires further analysis. In principle the number of photoexcited carriers should be proportional to the incident power. The observed quadratic dependence might be reflecting an increased absorption at the fixed laser energy, e.g. due to a GaAs gap reduction induced by a temperature increase.

This latter result points towards the physical origin of the observed Gaussian photon potential. The potential, schematized in Fig. 1b, represents the effective local energy of the optical cavity mode of an equivalent planar structure. It is this lateral gradient of the cavity mode energy that leads to light confinement and to new photonic modes (just like a quantum well leads to new confined electronic states defined by, though shifted from, the confinement potential). Similar ideas have been applied to model effective phonon potentials\(^{31}\). The resulting modes are red-shifted (with respect to the unperturbed planar structure) by the laser excitation, mapping out the shape of the focused laser intensity profile, reflecting that the photoexcited carriers induce a spatially dependent \textit{increase} of the refractive index of the GaAs spacer.

As mentioned in the introduction, either electronic or thermal phenomena can be at the basis of the observed phenomena. The different phenomena have however very different dynamics, typically reflecting pico- and nanosecond recombination when photoexcited car-
riers are responsible, and microsecond thermal diffusion times if heating is dominant. To clarify this issue, the dynamics of the photon potential was studied using modulated optical excitation. The modulation was achieved with an acousto-optic modulator, fed with a square wave generator of tunable frequency and duty cycle. The top left panel in Fig. 4 presents photoluminescence spectra as a function of the modulation frequency, for a duty cycle of 20%. From blue to red, the curves correspond to modulation frequencies ranging from 1 MHz to 1 Hz. The figure clearly shows the 3D confined spectra for slow modulation (equivalent to the cw experiments shown above), which evolve into the laterally unconfined planar situation with increasing modulation frequency. The top right panel in Fig. 4 displays the photoluminescence spectra variation with duty cycle for a square wave modulation of 400 kHz. This particular frequency was selected because it is where the discrete modes start to be noticed for the minimum duty cycle available (20%). From blue to red, the curves correspond to duty cycles ranging from 20% (0.5 µs on, 2.0 µs off) to 80% (2.0 µs on, 0.5 µs off), with steps of 5%. The times involved in the observed dynamics clearly point to a thermal phenomenon.

Assuming that the change in the refractive index that gives rise to the optical well is directly proportional to a local change in temperature due to the laser excitation, it is possible to describe the physical situation as a heat diffusion problem. A thermal diffusion model was used to obtain the final temperature in the spacer \( T(0, \infty) \) as a function of the modulation frequency and duty cycle (see the Supplemental Material for a detailed description of the calculations). In the bottom panel in Fig. 4 we show a comparison between the experimentally observed optical potential well depth (circles, derived from experimental curves as in the top left figure) and the theoretical results (red dashed line) obtained from the thermal diffusion model, as a function of the inverse of the modulation frequency \( f \). The data in the inset corresponds to the well depth as a function of the duty cycle time, obtained from the measurements shown in the top right panel, again compared with the theoretical curve derived from the thermal diffusion model. The thermal diffusivity \( \alpha_T = 10 \mu m^2/\mu s \) used to derive these curves was taken from the literature. Only one parameter needs to be adjusted in the model, namely \( z_0 \), the vertical position (positive towards the substrate and with origin at the cavity spacer) above which the sample maintains the room temperature value. We obtained \( z_0 = 7.5 \mu m \), which is a reasonable value, being a position located within the substrate and of the order of the DBR cavity vertical dimension.
The excellent agreement between experiment and theory shown in Fig. 4 clearly confirms thermal effects at the origin of the reported light trapping phenomena. Equivalent 3D confinement was obtained using a 514 nm laser, which is absorbed at the Ga$_{0.9}$Al$_{0.1}$As layers of the top DBR, thus only weakly coupling with the spacer layer. This implies that two-laser schemes can be implemented, with one laser used to seed the cavity emission, and the other to dynamically control the 3D light trapping. We note that the GaAs gap strongly depends on temperature above $\sim 100$ K$^{35}$. Below this temperature electronic nonlinearities (which are typically weaker but faster) should be dominant over the observed thermal effects. With both approaches rich photon potentials could be accessible using phase and intensity modulated illumination as that used in super-resolution microscopy and other photonic applications.

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Additional Information

In this supplemental material we provide a brief description of the experimental setup, and details on the theoretical models used to explain the spatial and spectral distribution of the optically confined modes, as well as their thermal origin.

I. EXPERIMENTAL SETUP

The experimental setup used, outlined in Fig. S1, consists of a 20x and 0.3NA microscope objective, a beam splitter and a set of two lenses: the collection lens, and an optional k-space-imaging lens. A Spectra Physics Ti-Sapphire continuous wave laser is used for the excitation (red dashed line in Fig. S1). The laser energy was set to 1.63 eV in order to excite carriers only in the GaAs spacer. The emitted photoluminescence (yellow) is collected by the same objective and imaged on the spectrometer entrance slit by the collection lens. The latter is movable in the direction perpendicular to this slit, enabling the acquisition of the spectra coming from different spatial positions.

The optional k-space lens can be positioned so that its focus coincides with the microscope objective back-focus. This technique, usually referred to as “k-space imaging”, is based on the
same ideas as that of reciprocal space imaging in transmission electron microscopy\textsuperscript{36}. Instead of imaging the real image of the sample, the collection lens images the plane corresponding to the microscope objective back-focus. All the light coming from the sample at a certain angle is focused on the same point on this plane. Therefore, the image formed at the spectrometer entrance slit has the different angular components of the emitted photoluminescence spatially separated, enabling the spectral analysis of the emission’s dispersion relation.

II. OPTICAL POTENTIAL WELL MODEL

![Graph showing confinement dependence of Bessel type modes and effective optical potential well.]

FIG. S2: Confinement dependence of Bessel type modes (black) and effective optical potential well (red). The solutions are given by the crossings between the curves. The red circles on the red curve are the energies of the measured modes. $P = 30\, \text{mW}$.

In the main text we presented the energy profile of the effective Gaussian potential generated by the laser excitation:

$$\varepsilon = \varepsilon_0 - \Delta \varepsilon_0 e^{-\frac{D_{\text{eff}}^2}{8\sigma^2}},$$

(3)

as well as the energy dependence of the confined Bessel modes:

$$\varepsilon = \sqrt{\varepsilon_0^2 + \frac{4\hbar^2 c^2}{n_{\text{cav}}^2} \frac{X_{ml}^2}{D_{\text{eff}}^2}}.$$  

(4)

We use Eq. (3) and Eq. (4) to calculate the energy of the photonic modes trapped in the Gaussian effective potential. Because of the symmetry of the problem not all Bessel modes are observed experimentally, so for the calculation we only use the orders with the right
symmetries. Since the order that corresponds to each energy is known, the dependence of ε with X_{ml} can be studied. This allows, making an approximation, to estimate ε_0 from Eq. (4). In Eq. (4) it is evident that ε depends on both the order of the mode (X_{ml}) and the effective diameter of the well (D_{eff}). In addition, D_{eff} also depends on the order of the mode. Therefore, D_{eff} and X_{ml} will depend on each other with some functionality such that when increasing X_{ml}, both D_{eff} and ε will increase. The simplest function that does just that is

\[ X_{ml} \propto D_{eff}^2. \]  

This seemingly arbitrary choice is, in fact, analogous to a second order Taylor expansion of X_{ml}, were the function X_{ml}(D_{eff}) continuous and derivable. Replacing Eq. (5) in Eq. (4), it follows that

\[ \varepsilon^2 = \varepsilon_0^2 + \alpha X_{ml}, \]  

being α a proportionality constant. This expression entails that applying a linear fit to \( \varepsilon^2(X_{ml}) \) we can determine \( \varepsilon_0 \).

![Graph showing the comparison of optical potential wells for three different excitation powers.](image)

**FIG. S3:** Comparison between the obtained optical potential wells for three different excitation powers.

Having estimated \( \varepsilon_0 \), D_{eff} can be calculated for each mode from Eq. (4). As we already mentioned, the solutions must comply with Eq. (4) and Eq. (3). Thus, knowing the effective diameter and energy of each mode, Eq. (3) can be used to obtain \( \Delta \varepsilon_0 \) and σ. This can be solved in a graphical form, where the energy curve for the well (Eq. (3)) must cross the energy curve of each optical mode (Eq. (4)) for a solution to the problem to exist. The
energy and effective diameter of each mode is given by the position of these crossings. As an illustration of the method, in Fig. S2 we present the result of the optical well potential model for an excitation power of 30 mW. The red curve represents the well, while the black curves show the dependence of the different modes with effective lateral confinement. The solutions are given by the crossings between the red and black curves. The experimental points are represented with the gray filled circles. An excellent agreement between model and experiment is observed for all the studied powers. In Fig. S3 the effective photon potentials obtained for three different excitation powers are compared. A deepening of the well is observed as the power is increased, while the width turns out to be constant, as expected from the invariant laser spot shape.

III. THERMAL DIFFUSION MODEL

FIG. S4: a) Radial temperature distribution as a function of time for a circular heat source of radius $R = 10 \, \mu m$ inside the GaAs spacer. The white dashed line marks the border of the heat source. b) Temperature distribution in the growth direction. The white dotted line marks the position of the GaAs spacer, i.e. the heat source. The latter is turned on at $t = 0 \, \mu s$ and off at $t \approx 0.36 \, \mu s$. 

17
Assuming that the change in the refractive index that gives rise to the optical effective potential is directly proportional to a local change in temperature due to the laser excitation, it is possible to describe the system as a heat diffusion problem:

\[
\frac{\partial T(r, t)}{\partial t} = \alpha_T \nabla^2 T(r, t) + Q(r, t),
\]

being \( T \) the temperature, \( \alpha_T \) the thermal diffusivity and \( Q \) a source term. Several approximations can be performed to provide a clear picture of the physics involved. Let us consider first the source term. Any temperature change must be due to the non-radiative relaxation of photo-excited carriers. Therefore, the temporal and spatial dependence of \( Q \) will be given by that of the e-h pairs excited by the incident laser. Because of the magnitude of the electronic gap, close to the laser excitation, the GaAs spacer is the only layer that will actually be excited by the 1.63 eV laser. Furthermore, the Al-rich layers on top and below the spacer act as energy barriers (because of the higher gap), so carriers will be confined in the spacer until they recombine. The heat source can therefore be considered not null only on the cavity spacer. In the lateral direction, the distribution will be given by the convolution of excitation profile (\( \sim 5 \mu m \) Gaussian spot) and the diffusion of the photoexcited carriers (i.e., a Gaussian distribution with a FWHM \( \approx 10 \mu m \) in our case).

Heat diffusion turns out to be highly anisotropic. Because the temperature gradients in the in-plane direction are considerably smaller than the one generated in the growth direction, lateral diffusion can be neglected. This qualitative argument was corroborated by a full finite-elements simulation considering the heat diffusion problem with an homogeneous heat source in a circular region of radius 10\( \mu m \) inside the GaAs spacer. Two color maps are presented in Fig. S4, showing the temperature distribution as a function of time in the radial (Fig. S4a) and growth (Fig. S4b) direction. It is clear that most of the heat diffuses towards the substrate, while the temperature profile in the radial direction changes mostly in amplitude, but not in shape. Therefore, the simpler one dimensional problem of the diffusion equation will be considered:

\[
\frac{\partial T(z, t)}{\partial t} = \alpha_T \frac{\partial^2 T(z, t)}{\partial z^2} + Q(z, t).
\]

Since the cavity spacer is very thin with respect to the size of the sample, the heat source will be approximated as

\[
Q(z, t) = \delta(z) \cdot q(t),
\]
where $\delta(z)$ is a Dirac delta. Regarding the temporal dependence, since the characteristic recombination times for photoexcited e-h pairs (that act here as an upper limit to the heating times) in these structures are around $7 \, \text{ns}$, and the modulation times we are working with are of the order of the $\mu$S, the heat generation can be considered instantaneous. That is, a photon will be instantaneously converted into heat at the spacer, and thus the time dependence of $Q$ will closely follow the time dependence of the laser excitation.

The considered problem has two distinct regimes: $q(t) = q$ during the ON-cycle, and $q(t) = 0$ during the OFF-cycle. Let’s consider the ON-cycle first. The solution to Eq. (8) has the following form:

$$T(z, t) = T_0 + \int_0^t \int_{-\infty}^{\infty} \frac{1}{\sqrt{4\pi \alpha_T(t-t')}} \exp\left(-\frac{(z-z')^2}{4\alpha_T(t-t')}\right) Q(z', t') dz' dt'$$

$$T(z, t) = T_0 + \int_0^t \frac{q}{\sqrt{4\pi \alpha_T(t-t')}} \exp\left(-\frac{(z-z')^2}{4\alpha_T(t-t')}\right) \delta(z') dz' dt'$$

$$T(z, t) = T_0 + \int_0^t \frac{q}{\sqrt{4\pi \alpha_T(t-t')}} \exp\left(-\frac{(z-z')^2}{4\alpha_T(t-t')}\right) dt'.$$

(10)

After the first ON-cycle, the temperature distribution will be given by

$$T(z, t_{\text{end}}) = T_0 + \int_0^{t_{\text{end}}} \frac{q}{\sqrt{4\pi \alpha_T(t_{\text{end}}-t')}} \exp\left(-\frac{(z-z')^2}{4\alpha_T(t_{\text{end}}-t')}\right) dt', \quad (11)$$

being $t_{\text{end}}$ the end time of the cycle. In the case when there is no heat source ($q(t) = 0$), but there is a non homogeneous temperature profile, the solution to Eq. (8) is:

$$T(z, t) = \int_{-\infty}^{\infty} \frac{T(z', t_{\text{end}})}{\sqrt{4\pi \alpha_T(t-t_{\text{end}})}} \exp\left(-\frac{(z-z')^2}{4\alpha_T(t-t_{\text{end}})}\right) dz'.$$

(12)

Comparing Eq. (10) and Eq. (12) it is evident that the latter is equal to the solution found for the problem of a heat source of the following form:

$$Q(z, t) = T(z, t) \cdot \delta(t - t_{\text{end}}).$$

(13)

Except for the very first ON-cycle, there is always a residual temperature distribution left from earlier cycles. To take this into account, we define the following effective heat source for each ON-cycle:

$$Q(z, t) = \delta(z) \cdot q \cdot \Theta(t - t_{0,i}) + T(z, t) \cdot \delta(t - t_{0,i}),$$

(14)
being \( t_{0,i} \) the initial time of ON-cycle \( i \) and \( \Theta(t_{0,i}) \) the Heaviside step function. The OFF-cycles, on the other hand, have the form

\[
Q(z, t) = T(z, t) \cdot \delta(t - t_{\text{end},i}),
\]

being \( t_{\text{end},i} \) the end time of the ON-cycle \( i \). The solution of the problem will then be

\[
T(z, t) = \int_{-\infty}^{\infty} \frac{T(z', t_{0,i})}{\sqrt{4\pi \alpha T(t - t_{0,i})}} e^{-\frac{(z - z')^2}{4\alpha T(t - t_{0,i})}} dz' + \int_{t_{0,i}}^{t} \frac{q}{\sqrt{4\pi \alpha T(t - t')}} e^{-\frac{z^2}{4\alpha T(t - t')}} dt',
\]

(16)
during the ON-cycles, and

\[
T(z, t) = \int_{-\infty}^{\infty} \frac{T(z', t_{\text{end},i})}{\sqrt{4\pi \alpha T(t - t_{\text{end},i})}} e^{-\frac{(z - z')^2}{4\alpha T(t - t_{\text{end},i})}} dz'.
\]

(17)
during the OFF-cycles. So far, we have not taken into account any heat sinks, therefore the system temperature would increase indefinitely, which is not physically correct. To address this, we impose a constant temperature plane at a distance \( z_0 \) from the spacer. To do so, we consider a negative image of the heat source at a distance \( 2z_0 \) from the real source\(^{37}\). The solution for the ON-cycles will then be

\[
T(z, t) = \int_{-\infty}^{\infty} \frac{T(z', t_{0,i})}{\sqrt{4\pi \alpha T(t - t_{0,i})}} e^{-\frac{(z - z')^2}{4\alpha T(t - t_{0,i})}} dz' + \int_{t_{0,i}}^{t} \frac{q}{\sqrt{4\pi \alpha T(t - t')}} \left[ e^{-\frac{z^2}{4\alpha T(t - t')}} - e^{-\frac{(z - 2z_0)^2}{4\alpha T(t - t')}} \right] dt'.
\]

(18)
The final temperature \( T(0, \infty) \) is, as already mentioned, proportional to the maximum refractive index change, and thus proportional to the optical potential well depth. From photoluminescence measurements performed as a function of temperature, we found that above 290 K and up to at least 340 K, the spectral position of the optical mode depends linearly with temperature with a slope of \( \frac{\partial \varepsilon}{\partial T} = 0.118 \pm 0.004 \text{ meV/ K} \). Therefore, a \( \sim 3.5 \text{ meV} \) optical well as the one presented in Fig. S2 corresponds to a radial difference en temperature of \( \sim 30 \text{ K} \).