We investigate the resonant cooling phenomena of a driven two-level radiator embedded in a photonic crystal structure. We find that cooling occurs even at laser-atom-frequency resonance. This happens due to the atomic dressed-states coupling with different strengths resulting in population redistribution and cooling. Furthermore, for off-resonant driving the two-level particle can be in either ground or excited bare-state, respectively. This may help in engineering of novel amplified optical devices as well as highly coherent light sources.

PACS numbers: 37.10.De, 42.50.Ct, 42.55.Tv

I. INTRODUCTION

The mechanical action of light in resonant interaction with atoms was investigated for instance in [1]. There, the laser cooling techniques were described in details. Various developments of those ideas were further implemented [2] (and references therein). Additionally, laser cooling effects of trapped ions in a standing- or a running-wave were studied in [3], while laser cooling of these systems via a squeezed environmental reservoir was investigated in Ref. [4]. Laser cooling of a trapped atom can be achieved in an optical cavity as well [4, 5]. Cavity-mediated laser cooling was further investigated in [6]. Dark-state laser cooling of a trapped ion using standing waves and cooling by heating in the quantum optics domain were recently discussed in [7] and [8], respectively. The experimental demonstration of ground-state laser cooling with electromagnetically induced transparency (EIT) [9] was performed in [10] for single-ion systems and in [11] for an ion chain. Furthermore, double-EIT ground-state laser cooling without blue-sideband heating was proposed in [12] while sideband cooling of atoms with the help of an auxiliary transition in [13], and Raman sideband cooling of a single atom in an optical tweezer in Ref. [14], respectively. Also, laser cooling of solids was addressed in [15] and [16], respectively.

Now, there is an increased interest to apply the above mentioned or related techniques to cool artificially created systems like nano-mechanical oscillators or quantum circuit systems [17]. In particular, schemes to ground-state cooling of mechanical resonators were proposed in [18]. Ground state cooling of a nanomechanical resonator in the nonresolved regime via quantum interference was the topic of Ref. [19]. Cooling a mechanical resonator via coupling to a tunable double quantum dot was investigated in [20], while ground-state cooling of a mechanical resonator coupled to two coupled quantum dots was proposed in [21]. The role of qubit dephasing in cooling a mechanical resonator by quantum back-action was discussed in [22]. There it was shown that ground-state cooling of a mechanical resonator can only be realized if the superconducting flux qubit dephasing rate is sufficiently low. A scheme for the cooling of a mechanical resonator via projective measurements on an auxiliary flux qubit which interacts with it was proposed in [23]. Furthermore, a flux qubit was experimentally cooled [24] using techniques somewhat related to the well-known optical sideband cooling methods. Finally, cooling a quantum circuit via coupling to a multiqubit ensemble was discussed in [25].

Here, we investigate the cooling dynamics of a two-level atom via a photonic crystal environment. The two-level emitter is trapped in such a medium, while the linear dimensions of the trapping area inside the material have to be of the order of several emission wavelengths or even more. The trapping geometry has to be carefully engineered in order to the Born-Markov approximation being valid. The trapped two-level emitter is pumped with a moderately intense coherent laser field that does not destroy the optical material due to insufficient applied intensities. When the corresponding generalized Rabi frequency is larger than the spontaneous decay rate, we have found an efficient cooling mechanism applicable even for resonant laser-atom interactions. The reason is various couplings of the involved atomic dressed-state transitions with the environmental reservoir leading to different decay rates that are also responsible for improving the cooling efficiency. While the two-level radiator is always in its lower dressed-state during cooling processes, in the bare-state frame it can be even in the excited state, i.e. the two-level atomic system is inverted. This situation is due to atom’s coupling with the photonic crystal material and does not occur in free-space. On the other side, in free space or modified reservoirs efficient cooling occurs when the generalized Rabi frequency is additionally of the order of the vibrational mode frequency. Note that quantum dynamics in photonic crystal environments was/is widely investigated (see, for instance, [26–30] and references therein) including various possibilities for trapping of atoms in such media [31, 32]. Moreover, laser cooling of a trapped particle in a harmonic potential with increased Rabi frequencies was investigated as well in [33].
The article is organized as follows. In Sec. II we describe the analytical approach and the system of interest. We apply here standard procedures, generalized to photonic crystals environments, to arrive at a master equation describing the atomic vibrational degrees of freedom only. Sec. III deals with the corresponding equations of motion and discussion of the obtained results. Both the steady-state and time-dependent quantum dynamics are analyzed there as well. Finally, the Summary is given in the last section, i.e. Sec. IV.

II. ANALYTICAL APPROACH

The Hamiltonian describing a two-level atomic system embedded in a photonic crystal environment and possessing an induced dipole $d$, frequency $\omega_0$, and interacting with a coherent source of frequency $\omega_L$, in a frame rotating at $\omega_L$, and in the dipole approximation is [34]:

$$
H = \sum_k \hbar \delta_k a_k^\dagger a_k + h\Delta S_z + h\Omega(S^+e^{i\mathbf{k}_L\cdot\mathbf{r}} + S^-e^{-i\mathbf{k}_L\cdot\mathbf{r}})
+ i\sum \{ \hat{g}_k \cdot \mathbf{d} \} (a_k^\dagger S^-e^{-i\mathbf{k}\cdot\mathbf{r}} - a_k S^+e^{i\mathbf{k}\cdot\mathbf{r}}).
$$

Here in the Hamiltonian (1), the first and the second terms describe, respectively, the free environmental electromagnetic field (EMF) vacuum modes with $\delta_k = \omega_k - \omega_L$, as well as the free atomic energy with $\Delta = \omega_0 - \omega_L$. The third one characterizes the interaction of the two-level emitter localized at position $\mathbf{r}$ with the external coherent laser field, while $\Omega$ being the corresponding Rabi frequency. The last term considers the qubit’s interaction with environmental EMF vacuum modes reservoir with $g_k$ describing the corresponding interaction strength. The atomic bare-state operators $S^\pm = |2\rangle\langle 1| + |1\rangle\langle 2|$ obey the commutation relations for su(2) algebra: $[S^+ , S^-] = 2S_z$ and $[S_z , S^\pm] = \pm S^\pm$. Here, $S_z = (\langle 2|2| - |1\rangle\langle 1|)/2$ is the bare-state inversion operator. $|2\rangle$ and $|1\rangle$ are the excited and ground state of the qubit, respectively, while $a_k$ and $a_k^\dagger$ are the creation and the annihilation operators of the EMF, and satisfy the standard bosonic commutation relations, i.e., $[a_k , a_k^\dagger] = \delta_{kk}$, and $[a_k , a_k^\dagger] = [a_k^\dagger , a_k^\dagger] = 0$.

We are interested in the laser dominated regime and transfer our investigations in the laser-dressed picture: $|2\rangle = \cos \theta |2\rangle - \sin \theta |1\rangle$ and $|1\rangle = \sin \theta |2\rangle + \cos \theta |1\rangle$ with cot $2\theta = \Delta/(2\Omega)$. Further, the quantized vibrations of the atom’s center-of-mass motion are represented as usual in the Lamb-Dicke limit, i.e. $k_L = \eta (b^\dagger b - 1)$, where $b^\dagger (b)$ satisfy the single-mode bosonic commutation relations [35]. Expanding the laser-atom interaction Hamiltonian to the second order in the small parameter $\eta$ as well as eliminating the EMF operators in the standard way and in the Born-Markov approximations [34] one arrives then at the following dressed-state master equation:

$$
\frac{d\rho}{dt} = -i[H_0, \rho] - \frac{\gamma_o}{4}\sin^2 2\theta \{ R_z^2 \rho - 2R_z \rho R_z + \rho R_z^2 \}
- \gamma_+ \cos^4 \theta \{ R^+ R^- \rho - 2R^+ \rho R^- + \rho R^+ R^- \}
- \gamma_- \sin^4 \theta \{ R^- R^+ \rho - 2R^- \rho R^+ + \rho R^- R^+ \},
$$

where we have set $\hbar = 1$ and

$$
H_0 = \nu b^\dagger b + \bar{\Omega} R_z + i\eta \Omega (R^+ - R^-)(b^\dagger + b).
$$

Here, the vibrational mode frequency is given by $\nu$ with $\nu b^\dagger b$ being its free energy. An additional second term, proportional to $\eta^2$, in Eq. (3) leads to a shift of the oscillator frequency as well as to contributions proportional to higher orders in $\eta$ and, thus, was not taken into account here. Further, the atomic dressed-basis operators are defined as: $R_z = |2\rangle\langle 2| - |1\rangle\langle 1|$, $R^+ = |2\rangle\langle 1|$, and $R^- = |1\rangle\langle 2|$, while obeying the commutation relations $[R^+, R^-] = R_z$, and $[R_z, R^\pm] = \pm 2R^\pm$, respectively. $\gamma_\pm$ and $\gamma_o$ are the single-atom dressed-state spontaneous decay rates at the frequencies $\omega_L \pm 2\bar{\Omega}$ and $\omega_L$ [35], respectively, while [3]

$$
\bar{\rho} = \frac{1}{2} \int_{-1}^{1} dx w(x)e^{i\eta(b^\dagger b - 1)x} \rho e^{-i\eta(b^\dagger b - 1)x}.
$$

The angular distribution of the spontaneous emission is characterized by the function: $w(x) = \frac{1}{2}(1 + x^2)$. Notice that in free space environments $\gamma_o = \gamma_\pm$. This is not more the case if the atom is surrounded by a modified environmental electromagnetic field reservoir such as photonic crystals. Depending on density of modes distributions around the relevant dressed-state transitions these decay rates can differ substantially leading to inversion in the bare-states [27]. Finally, in Eq. (2), one have performed the secular approximation, i.e. we have neglected the terms oscillating at the generalized Rabi frequency $\Omega = \sqrt{\Omega^2 + (\Delta/2)^2}$ or higher frequencies. This approximation is valid for $\Omega \gg \gamma_o, \pm$, i.e. the spectral bands of the Mollow spectrum [36] are well distinguished.

In what follows, we shall apply the standard procedure to eliminate the atomic degrees of freedom from the Eq. (2). This is valid when the atomic subsystem is faster than the vibrational ones. We proceed as follows: i) expand the exponents in Eq. (4) to the second order ii) go in a rotating frame at the generalized Rabi frequency $\Omega$; and iii) take the trace over atomic degrees of freedom. Then one obtain the following master equation for the vibrational mode that still contains the atomic operators:

$$
\frac{d\rho}{dt} = \eta \Omega ([B(t), \rho_{21}] e^{2i\bar{\Omega} t} - [B(t), \rho_{21}] e^{-2i\bar{\Omega} t})
- \alpha \eta^2 \{ (\gamma_+ \cos^4 \theta + \frac{\gamma_o}{4}\sin^2 2\theta) [B(t), B(t) \rho_{22}]
+ (\gamma_- \sin^4 \theta + \frac{\gamma_o}{4}\sin^2 2\theta) [B(t), B(t) \rho_{11}] + H.c. \},
$$

(5)
FIG. 1: (color online) The steady-states of (a) the mean-photon number $n = \langle b^\dagger b \rangle_s$ and (b) dressed-state $\langle R_s \rangle_s$ (solid line) as well as bare-state $2\langle S_s \rangle_s$ (long-dashed curve) inversion operator versus normalized detuning. Here, $\eta = 0.1$ whereas $\Omega/\gamma = 5$ with $\gamma_+ = \gamma_- = \gamma_0 \equiv \gamma$. In (a) the solid, long-dashed and short-dashed lines are for $\nu/\gamma = 2$, 6 and 12, respectively.

Here $\alpha = 2/5$ and $B(t) = be^{-i\nu t} + b^\dagger e^{i\nu t}$ while $\rho_{\alpha\beta} = \langle \alpha | \rho_0 | \beta \rangle$ with $\{\alpha, \beta \in \{1, 2\}\}$, and $\rho_0$ being the atomic density matrix operator only. To first order in $\eta$, from Eq. (2), one can obtain:

$$\rho_{21}(t) = \eta \Omega \{ B(t) \rho_{11} - \rho_{22} B(t) \} e^{2i\Omega t},$$

(6)

where $\rho_{12}(t) = (\rho_{21}(t))^\dagger$, and

$$B(t) = \frac{be^{-i\nu t}}{\Gamma_\perp + i(2\Omega - \nu)} + \frac{b^\dagger e^{i\nu t}}{\Gamma_\perp + i(2\Omega + \nu)},$$

with $\Gamma_\perp = \gamma_0 \sin^2 2\theta + \gamma_+ \cos^4 \theta + \gamma_- \sin^4 \theta$. Inserting Eq. (3) in Eq. (5) and keeping only resonant contributions one arrive at the following master equation describing the vibrational mode alone:

$$\frac{d}{dt}\rho = -A_{s}^* [b^\dagger, b]\rho - A_{s}^* [b, b^\dagger \rho] + H.c.,$$

(7)

where “*” means complex conjugation and

$$A_{s}^* = \Gamma_0 + \frac{\eta \Omega^2 \langle R_{11} \rangle_s}{\Gamma_\perp + i(2\Omega - \nu)}; \quad A_{s}^* = \Gamma_0 - \frac{\eta \Omega^2 \langle R_{22} \rangle_s}{\Gamma_\perp + i(2\Omega - \nu)};$$

with $\Gamma_0 = \alpha \eta^2 (\gamma_- \sin^4 \theta \langle R_{11} \rangle_s + \gamma_+ \cos^4 \theta \langle R_{22} \rangle_s + \gamma_0 \sin^2 2\theta)$. The steady-state values of the atomic operators $\langle R_{os} \rangle_s = \langle \alpha | \rho_0 | \alpha \rangle_s$, $\alpha = \{1, 2\}$, are obtained from Eq. (2) without taking into account the vibrational degrees of freedom, namely [32]:

$$\langle R_{11} \rangle_s = \frac{\gamma_+ \cos^4 \theta}{\gamma_+ \cos^4 \theta + \gamma_- \sin^4 \theta};$$

$$\langle R_{22} \rangle_s = 1 - \langle R_{11} \rangle_s.$$  

(8)

That is, we have considered that the vibrational degrees of freedom do not modify the population distribution among the involved dressed-states that are governed by the external applied coherent laser field.

In the next section, with the help of Eq. (7), we shall describe the cooling processes of a pumped two-level emitter embedded in a photonic crystal.

In the next section, with the help of Eq. (7), we shall describe the cooling processes of a pumped two-level emitter embedded in a photonic crystal.

FIG. 2: (color online) Same as in Fig. 1 with $\Omega/\gamma = 5$, $\gamma_-/\gamma_+ = \gamma_0/\gamma_+ = 0.2$. In (a) the solid, long-dashed and short-dashed curves are for $\nu/\gamma = 2$, 6 and 12, respectively.

III. COOLING PHENOMENON

In the Heisenberg picture, the master equation (7) can be represented as follows:

$$\frac{d}{dt} \langle Q \rangle = -A_{s}^* \langle [Q, b^\dagger b] \rangle - A_{s}^* \langle [Q, b^\dagger b^\dagger] \rangle + H.c.,$$

(9)

where $Q$ is any vibrational operator. Note that, in general, for the non-Hermitian operators $Q$, the H.c. terms should be evaluated without conjugating $Q$, i.e., by replacing $Q^\dagger$ with $Q$ in the Hermitian conjugate parts. With the help of Eq. (9), one can easily obtain the quantum dynamics of the mean phonon number in the vibrational mode, that is:

$$\frac{d}{dt} \langle b^\dagger b \rangle = -(A_{s}^- - A_{s}^+) \langle b^\dagger b \rangle + A_{s}^+, \quad (10)$$

where $A_{s}^- = A_{s} - A_{s}^+$, while $A_{s}^+ = A_{s} + A_{s}^*$. Evidently, cooling occurs for $A_{s}^- > A_{s}^+$. Particularly, the cooling rate $C = A_{s}^- - A_{s}^+$ is:

$$C = -\frac{2(\eta \Omega)^2 \Gamma_\perp \langle R_{22} \rangle_s}{\Gamma_\perp^2 + (2\Omega - \nu)^2},$$

(11)

where the steady-state dressed-state inversion operator is given by $\langle R_{s} \rangle_s = \langle R_{22} \rangle_s - \langle R_{11} \rangle_s$. One can observe that cooling occurs always when the two-level emitter is in the lower dressed-state, i.e. $\langle R_{2} \rangle < 0$. In the steady-state, we have from Eq. (11) and Eq. (11) that $\langle b^\dagger b \rangle_s = A_{s}^+ / C$ or:

$$\langle b^\dagger b \rangle_s = \frac{\langle R_{22} \rangle_s}{\langle R_{11} \rangle_s - \langle R_{22} \rangle_s} + \frac{\Gamma_0 (\Gamma_\perp^2 + (2\Omega - \nu)^2)}{(\eta \Omega)^2 \Gamma_\perp^2 (\langle R_{11} \rangle_s - \langle R_{22} \rangle_s)}.$$  

(12)

In the next subsections, we shall rigorously analyze the steady-state as well as the time-dependent quantum dynamics of the cooling effect characterizing the particular system described here.

A. The steady-state cooling dynamics

Fig. (1b) shows the mean-number of vibrational quanta according to Eq. (12) and for standard situations, that
is, when the pumped two-level emitter interacts with the usual vacuum modes of the environmental electromagnetic field reservoir [3, 32]. Cooling efficiency improves, i.e. \( n \ll 1 \), if the generalized Rabi frequency \( 2\Omega \) approaches the vibrational frequency \( \nu \) while the external pumping coherent field is evidently off-resonant. As it was already mentioned, cooling always occurs when the two-level atom is in its ground dressed-state, and for free space, the atom is also in its ground bare-state, respectively, (see Fig. 1b) where the bare-state inversion operator as a function of \( \gamma_-/\gamma_+ \). Here, \( \Delta = 0 \) and \( \gamma_0 \equiv \gamma_- \) while all other parameters are the same as in Fig. 4.

![Fig. 4](image)

**FIG. 4:** (color online) The steady-states of (a) the mean-phonon number \( n = \langle b^\dagger b \rangle_s \) and (b) dressed-state \( \langle R_s \rangle_s \) (solid line) as well as bare-state \( \langle S_s \rangle_s \) (long-dashed curve) inversion operator as a function of \( \gamma_-/\gamma_+ \). Here, \( \Delta/(2\Omega) = -0.5 \) while all other parameters are the same as in Fig. 4.

In this respect, Fig. (3a) depicts the mean-vibrational-phonon number when the two-level emitter is embedded in a modified environmental electromagnetic field reservoir such as photonic crystals. Interestingly, cooling occurs even at laser-atom resonance. This happens because the decay rates at various involved dressed-state transitions can differ substantially which is not the case in free space at resonance. Here, again, the atom is in its lower dressed-state while the bare-state inversion is zero meaning that the atom is equally distributed on the two bare states (see Fig. 3b). Furthermore, in Fig. 4, we show the cooling processes when the two-level emitter can be even in the excited bare-state. This may be of particular interest for engineering of highly coherent laser sources, for instance, where the vibrational degrees of freedom do not influence the photon statistics. Efficient cooling via modified environmental electromagnetic field reservoirs occurs also for positive detunings although the atom will be in its lower dressed/bare-state, respectively. Thus, cooling takes place for both positive or negative detunings. This feature is not proper for free space setups, i.e. when the two-level atom is surrounded by the usual vacuum modes of the environmental electromagnetic field reservoir. The explanation is as follows: the laser-dressed atom decays on transition \( |2 \rangle \rightarrow |1 \rangle \) with a decay rate \( \gamma_+ \cos^4 \theta \), while on \( |1 \rangle \rightarrow |2 \rangle \) transition with \( \gamma_- \sin^4 \theta \), respectively. Cooling occurs when the two-level emitter is in its lower dressed-state, \( |1 \rangle \), because in this state the phonon generation processes are minimized, and, therefore, \( \gamma_+ \cos^4 \theta \) should be larger than \( \gamma_- \sin^4 \theta \). In free space \( \gamma_+ = \gamma_- \) and, thus, cooling is achieved for positive detunings only since in this case \( \cos^4 \theta > \sin^4 \theta \). Note that \( \cos^2 \theta = (1 + \Delta/(2\Omega))/2 \) while \( \sin^2 \theta = (1 - \Delta/(2\Omega))/2 \). However, due to coupling of the dressed-atom with photonic crystal environments, cooling can take place for negative detunings as well and still \( \gamma_+ \cos^4 \theta > \gamma_- \sin^4 \theta \), i.e. when \( \gamma_+ > \gamma_- \).

**B. The time-dependent cooling dynamics**

The time-dependent quantum dynamics for the mean values of the dressed-state atomic inversion operator, dressed-state coherences, as well as the mean-phonon number of vibrational quanta is described by the following expressions (see Eq. 2 and Eq. 11):

\[
\langle R_s(t) \rangle = \langle R_s(0) \rangle e^{-2\gamma_+t} + \langle R_s \rangle_s,
\]

\[
\langle R^+_s(t) \rangle = \langle R^+_s(0) \rangle e^{-\Gamma_-t}, \quad \text{with} \quad \langle R^- \rangle = \langle R^+ \rangle^\dagger,
\]

\[
\langle b^\dagger b \rangle_t = \langle b^\dagger b \rangle_0 - \langle b^\dagger b \rangle_s e^{-\Gamma t} + \langle b^\dagger b \rangle_s,
\]

where \( \gamma_s = \gamma_+ \cos^4 \theta + \gamma_- \sin^4 \theta \), while \( \langle R_s(0) \rangle = \cos^2 \theta \langle S(0) \rangle + \sin^2 \theta \langle S(0) \rangle - \sin^2 \theta \langle S(0) \rangle + \langle b^\dagger b \rangle_0 \) and \( \langle b^\dagger b \rangle_0 \) are the initial conditions for the mean values of atomic inversion operator, coherences, and vibrational phonon number, respectively. Note that the expressions (13) are valid for \( t \gg (2\Omega)^{-1} \) and \( \eta \Omega < \max\{\gamma_+, \gamma_0\} \ll 2\Omega \). Furthermore, our cooling approach requires that \( (2\gamma_0, \Gamma_-) \gg C \) which, in principle, can always be arranged. Particularly, for \( \{\gamma_0, (2\Omega - \nu) \approx \gamma_+ \) and \( \langle R_s \rangle_s \approx -1 \) one has \( C \approx (\eta \Omega/\gamma_+)^2 \gamma_+. \) When \( \Omega/\gamma_+ = 5 \) and \( \eta = 0.1 \), cooling occurs for \( t \gg 4/\gamma_+ \). For typical values of spontaneous decay rates at optical frequencies, \( \gamma_+ \sim 10^8 - 10^9 \text{Hz} \), cooling is achieved in microseconds.
IV. SUMMARY

In summary, we have investigated the cooling efficiency of a laser-pumped two-level emitter placed in a modified surrounding electromagnetic field reservoir like photonic crystals. Particularly, cooling occurs for positive or negative laser-atom detunings as well as at resonance. Furthermore, the two-state particle can be even in the excited bare-state during cooling processes facilitating sensitive applications towards better coherent light sources as well as light amplification.

Acknowledgments

G.-x. Li is grateful to the financial support from National Natural Science Foundation of China (Grant No. 61275123) and the National Basic Research Program of China (Grant No. 2012CB921602).

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