Electromagnetic radiation amplification by means of a driven two-level system damped by broadband squeezed vacuum reservoir

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Abstract. It is shown that a two-level quantum system with broken inversion symmetry possessing dipole moment operator with permanent non-equal diagonal matrix elements, driven by external semiclassical monochromatic high-frequency electromagnetic (laser) field and damped by broadband squeezed vacuum reservoir can amplify or absorb weak probe electromagnetic radiation field of much lower frequency depending on the reference phase of the squeezed field. Varying the degree of squeezing allows to alter the magnitude of the absorption or amplification rate.

1. Introduction

It has already been shown recently [1] that a two-level quantum system (also known as two-level atom) possessing dipole moment operator with permanent non-equal diagonal matrix elements and driven by external semiclassical monochromatic high-frequency electromagnetic (EM) field can amplify weak probe EM radiation field at much lower frequency. This important property of a relatively simple quantum system may hold promise for developing practically feasible techniques for the manipulation of the EM radiation in the terahertz range of frequencies. Terahertz radiation occupies a range of frequencies between microwaves and infrared waves, and is loosely defined to extend from about 0.3 THz (a wavelength of 1 mm) to around 30 THz (a wavelength of 10 µm). Terahertz waves are promising for applications in numerous fields of applied science and technologies because they are nonionizing, noninvasive, very sensitive to various kinds of resonances and penetrable to many materials. Nevertheless, the insufficient development of viable technics for the generation, amplification, modulation and any other kind of manipulation for the terahertz radiation range represents major impediment to broad usage of the aforesaid advantages in common everyday practice despite obvious successes in the field of the terahertz radiation generation achieved so far [2]. It seems highly desirable to employ the simplest of available physical systems as the means for the EM radiation manipulation in the terahertz range. To our knowledge, the idea that a two-level quantum system driven by high-frequency laser field can generate EM field of much lower frequency if its dipole operator possesses permanent non-equal diagonal matrix elements in contrast to conventional assumption that only the non-diagonal matrix elements persist, was put forward in [3]. This effect was further studied in [1, 4, 5], where the properties of this low-frequency radiation were thoroughly investigated.
for the case of a two-level system driven by external EM field and simultaneously interacting with a dissipative vacuum reservoir. This research is focused on furthering the investigation of the above-mentioned phenomenon of the EM radiation amplification and absorption to the case of an externally driven two-level system with broken inversion symmetry interacting with a squeezed vacuum dissipative reservoir, which properties can be tuned appropriately in order to control the degree of the absorption or amplification and/or to switch between these two modes of operation.

2. Model Hamiltonian

In this study we consider a two-level atom with ground state \( |g\rangle \), excited state \( |e\rangle \), transition frequency \( \omega_0 \) and the electric dipole moment \( \mathbf{d} \), driven by external classical monochromatic field \( \mathbf{E}(t) = \mathbf{E} \cos(\omega_f t) \) with the amplitude \( \mathbf{E} \) and the frequency \( \omega_f \). The atom is also coupled to a reservoir \( B \) made of a plurality of modes of quantized electromagnetic field being in the squeezed vacuum state. It is assumed that the frequency Lamb shift due to interaction with the reservoir is already incorporated into the atomic transition frequency \( \omega_0 \). Thus, the model Hamiltonian reads

\[
H = H_S(t) + \hbar \sum_k \omega_k b^\dagger (\omega_k) b(\omega_k) + \sum_k (g(\omega_k) S^+ b(\omega_k) + g^*(\omega_k) S^- b^\dagger (\omega_k)).
\]

(1)

Here \( S^+ = |e\rangle\langle g| \) and \( S^- = |g\rangle\langle e| \) are the usual raising and lowering atomic operators and \( S^z = \frac{1}{2}(|e\rangle\langle e| - |g\rangle\langle g|) \) is the atomic population inversion operator. The term

\[
H_S(t) = \hbar \omega_0 S^z + \frac{\hbar}{2} (e^{i\omega_f t} + e^{-i\omega_f t}) \left[ \Omega_R S^+ + \Omega_b S^- - \mathbf{E}_{dg} |e\rangle\langle e| - \mathbf{E}_{dd} |g\rangle\langle g| \right] ,
\]

(2)
describes the interaction between the driving field and the atom. Here \( \Omega_R = -\mathbf{E}_{dg}/\hbar \) is the Rabi frequency being made real and positive by the appropriate choice of the phase factors of the states \( |e\rangle \) and \( |g\rangle \), and \( \mathbf{d}_{eg} = e\langle e|\mathbf{r}|g\rangle \), \( \mathbf{d}_{ge} = e\langle g|\mathbf{r}|e\rangle \), \( \mathbf{d}_{ee} = e\langle e|\mathbf{r}|e\rangle \), \( \mathbf{d}_{gg} = e\langle g|\mathbf{r}|g\rangle \) are the atomic dipole moment operator matrix elements. As a rule, it is assumed that \( \mathbf{d}_{ee} = \mathbf{d}_{gg} = 0 \), because typical physical systems, like atoms and molecules, possess the inversion symmetry, and each of the states \( |g\rangle \) and \( |e\rangle \) is either symmetric or antisymmetric. Contrary to this view, we assume that the inversion symmetry of the system in question is violated, \( \mathbf{d}_{ee} \neq \mathbf{d}_{gg} \), and the interaction Hamiltonian (2) can be identically transformed into

\[
H_S(t) = \hbar \omega_0 S^z + \frac{\hbar}{2} (e^{i\omega_f t} + e^{-i\omega_f t}) \left[ \Omega_R S^+ + \Omega_b S^- + \delta_a S^z - \frac{\delta_s}{2} (|e\rangle\langle e| + |g\rangle\langle g|) \right] \]

(3)

with \( \delta_a = \mathbf{E}(\mathbf{d}_{gg} - \mathbf{d}_{ee})/\hbar \) and \( \delta_s = \mathbf{E}(\mathbf{d}_{gg} + \mathbf{d}_{ee})/\hbar \). The term proportional to \( \delta_a \) does not influence the dynamics of the system and can be omitted, while the term proportional to the symmetry violation parameter \( \delta_s \) is retained. A broadband squeezed vacuum reservoir is characterized by the density operator [6]

\[
\rho_{\text{vac},sq}^{\text{vac}} = \hat{S}(\xi)|\text{vac}\rangle\langle \text{vac}|\hat{S}^\dagger (\xi), \quad \hat{S}(\xi) = \prod_k \exp \left( \xi^* b(\Omega_p - \omega_k)b(\Omega_p + \omega_k) - h.c. \right) ,
\]

(4)

where \( \xi = r \exp(i\theta) \), with \( r \) being the squeeze parameter and \( \theta \) being the reference phase for the squeezed field.
3. Equations of Motion for Atomic Variables

The approximate master equation for the atomic reduced density operator $\rho_S(t)$ in the Schrödinger picture is

$$\frac{\partial \rho_S(t)}{\partial t} = -i\hbar\{H_S(t), \rho_S(t)\} - \frac{1}{2}\Gamma(N(r) + 1)(S^+ S^- \rho_S(t) + \rho_S(t) S^- S^+) - \frac{1}{2}\Gamma N(r)(S^- S^+ \rho_S(t) + \rho_S(t) S^+ S^-) - \Gamma (M^*(r, \theta)e^{i2\Omega \rho t} S^- \rho_S(t) S^+ + M(r, \theta)e^{-i2\Omega \rho t} S^+ \rho_S(t) S^-),$$

where $\Gamma$ is the radiative damping constant, $M(r, \theta) = \cosh(r) \sinh(r) \exp(i\theta)$, $N(r) = \sinh^2(r)$. This equation is derived assuming that the system-reservoir and the system-field interactions are weak, and the reservoir correlation time is small compared with the time $t$ of observation.

In what follows, it is assumed that the resonant frequency of the squeezing device $\Omega_p$ coincides with the atomic transition frequency $\omega_0$. So, a closed set of equations follows from Eq. (5):

$$\frac{d\langle \tilde{S}^-(t) \rangle}{dt} = -\left(\frac{\Gamma}{2}(2N(r) + 1) + i\Delta + i\frac{\delta_0}{2}(e^{i\omega_f t} + e^{-i\omega_f t})\right) \times \langle \tilde{S}^-(t) \rangle + \Gamma M(r, \theta)\langle \tilde{S}^+(t) \rangle e^{i2(\omega_f - \omega_0)t} + \Gamma R(1 + e^{-i2\omega_f t})\langle \tilde{S}^z(t) \rangle,$$

$$\frac{d\langle \tilde{S}^+(t) \rangle}{dt} = -\left(\frac{\Gamma}{2}(2N(r) + 1) - i\Delta - i\frac{\delta_0}{2}(e^{i\omega_f t} + e^{-i\omega_f t})\right) \times \langle \tilde{S}^+(t) \rangle + \Gamma M^*(r, \theta)\langle \tilde{S}^-(t) \rangle e^{-i2(\omega_f - \omega_0)t} + \Gamma R(1 + e^{-i2\omega_f t})\langle \tilde{S}^z(t) \rangle,$$

$$\frac{d\langle \tilde{S}^z(t) \rangle}{dt} = -\frac{\Omega_R}{2}(e^{i\omega_f t} + e^{-i\omega_f t})\langle \tilde{S}^+(t) \rangle e^{i\omega_f t} + \langle \tilde{S}^-(t) \rangle e^{-i\omega_f t},$$

where $\langle \tilde{S}^\pm(t) \rangle = \pm i\langle S^\pm(t)e^{\mp i\omega_f t} \rangle$ are slowly varying parts of the atomic operators and $\Delta = \omega_0 - \omega_f$ is the detuning between the atomic transition frequency and the frequency of the driving field. The system of equations (6-8) can be solved numerically for the case of the resonant driving field $\omega_f = \omega_0$, $\Delta = 0$ by means of the technique employed earlier in [7], where the components of the vector $\tilde{X}(t) = (\langle \tilde{S}^-(t) \rangle, \langle \tilde{S}^+(t) \rangle, \langle \tilde{S}^z(t) \rangle)$ are decomposed as

$$X_i(t) = \sum_{l=-\infty}^{+\infty} X_i^{(l)}(t)e^{il\omega_f t}, \quad i = 1, 2, 3,$$

and the slowly varying amplitudes $X_i^{(l)}(t)$ obey the system of equations

$$\frac{d}{dt} X_1^{(l)}(t) = \left(\frac{\Gamma}{2}(2N(r) + 1) + il\omega_f\right) X_1^{(l)}(t) - \frac{i\delta_0}{2}(X_1^{(l-1)}(t) + X_1^{(l+1)}(t)) + \Gamma M(r, \theta) X_2^{(l)}(t) + \Gamma R(X_3^{(l-2)}(t) + X_3^{(l)}(t)),$$

$$\frac{d}{dt} X_2^{(l)}(t) = \left(\frac{\Gamma}{2}(2N(r) + 1) + il\omega_f\right) X_2^{(l)}(t) + \frac{i\delta_0}{2}(X_2^{(l-1)}(t) + X_2^{(l+1)}(t)) + \Gamma M^*(r, \theta) X_1^{(l)}(t) + \Gamma R(X_3^{(l-2)}(t) + X_3^{(l)}(t)),$$

$$\frac{d}{dt} X_3^{(l)}(t) = -\frac{\Omega_R}{2}(X_2^{(l-2)}(t) + X_2^{(l)}(t) + X_1^{(l)}(t) + X_1^{(l+2)}(t)).$$
4. Probe Field Absorption-Amplification Spectrum

To study absorption and amplification of weak probe EM radiation by the system (1), this system is to be illuminated simultaneously by a weak probe field of frequency \( \omega_p \), amplitude \( E_p \) and arbitrary phase \( \phi_p \). The effect of this field can be accounted for by an additional term

\[
H_{SP}(t) = -eE_p \hat{r} \cos(\omega t + \phi_p).
\]

Here we assume that the probe field is linearly polarized along the same direction as the driving field. The term (12) can be reshaped in full analogy with the term (2) as

\[
H_{SP}(t) = \frac{\hbar}{2} \left( e^{i(\omega t + \phi_p)} + e^{-i(\omega t + \phi_p)} \right) \left[ \Omega_R^{(p)} (S^+ + S^-) + \delta_{a}^{(p)} S^z \right],
\]

where \( \Omega_R^{(p)} = -E_p d_{eg}/\hbar \) is the probe field Rabi frequency, which is already real due to the proper choice of the arbitrary phase factors of the states \( |e\rangle \) and \( |g\rangle \) in the expression for \( \Omega_R \) before, and can be made positive by the proper choice of the phase \( \phi_p \), and \( \delta_{a}^{(p)} = E_p (d_{gg} - d_{ee})/\hbar \) is the correspondent symmetry violation parameter. It is assumed that the probe field is much weaker than the driving field, so that it does not affect significantly the evolution of the system under the influence of the driving field. Then, the steady-state absorption spectrum of the probe field is given by the Fourier transform of the two-time commutators of the atomic operators as [6]

\[
W_a(\omega) = 2\hbar \omega \text{Re} \sum_{\lambda,\beta} \epsilon_\lambda(\omega) \epsilon_\beta(-\omega) \int_0^\infty \lim_{t \to \infty} dt e^{i\omega t} \left\langle \left[ \hat{Q}_\lambda(\omega, t + \tau), \hat{Q}_\beta(-\omega, t) \right] \right\rangle,
\]

where \( \epsilon_\lambda(-\omega) = \epsilon_\lambda^*(\omega) \), \( \hat{Q}_\lambda(-\omega) = \hat{Q}_\lambda^*(\omega) \) in general, and for the particular case of the probe field term in the form (12)

\[
\epsilon_1(\omega) = \frac{\Omega_R^{(p)}}{2} e^{i\phi_p}, \quad \epsilon_2(\omega) = \frac{\delta_a^{(p)}}{2} e^{i\phi_p}, \quad \hat{Q}_1(\omega) = S^+ + S^-, \quad \hat{Q}_2(\omega) = S^z, \quad \lambda, \beta = 1, 2.
\]

This function \( W_a(\omega) \) determines the rate at which the system, monochromatically driven by the strong EM field, absorbs energy from the weak probe field in terms of the two-time correlation functions of the system operators. Hence, all the commutators in (14) are to be calculated in the absence of the probe field and presence of the driving field. Variation of \( W_a(\omega) \) as a function of the probe field frequency \( \omega \) gives the absorption spectrum of the field. If we denote

\[
\tilde{S}_1(t) = \tilde{S}^-(t), \quad \tilde{S}_2(t) = \tilde{S}^+(t), \quad \tilde{S}_3(t) = \tilde{S}^z(t),
\]

then, according to the so-called quantum regression hypothesis [8], the commutator correlation functions

\[
D_{ij}(t, t + \tau) = \langle [\tilde{S}_i(t + \tau), \tilde{S}_j(t)] \rangle, \quad i, j = 1, 2, 3.
\]

satisfy almost the same set of equations of motion (6-8) as the correspondent averages \( \langle \tilde{S}^-(\tau) \rangle \), \( \langle \tilde{S}^+(\tau) \rangle \) and \( \langle \tilde{S}^z(\tau) \rangle \):

\[
\frac{dD_{ij}(t, t + \tau)}{d\tau} = -\left( \frac{\Gamma}{2} (2N(r) + 1) + i\Delta + \frac{\delta_a}{2} (e^{i\omega f(t+\tau)} + e^{-i\omega f(t+\tau)}) \right) D_{ij}(t, t + \tau) +
\]

\[+\Gamma M(r, \theta)D_{2j}(t, t + \tau)e^{ij(\omega_f-\omega_0)t} + \Omega_R(e^{i2\omega_f(t+\tau)} + 1)D_{3j}(t, t + \tau),
\]

\[\text{(18)}\]
\[
\begin{align*}
\frac{dD_{2j}(t, t + \tau)}{d\tau} &= - \left( \frac{\Gamma}{2}(2N(r) + 1) - i\Delta \right) - i\frac{\delta_\alpha}{2}(e^{i\omega_f(t+\tau)} + e^{-i\omega_f(t+\tau)}) D_{2j}(t, t + \tau) + \\
&\quad + \Gamma M^*(r, \theta) D_{1j}(t, t + \tau)e^{-i(\omega_f - \omega_0)t} + \Omega_R(1 + e^{-i2\omega_f(t+\tau)}) D_{3j}(t, t + \tau), \\
&\quad \quad \tag{19}
\end{align*}
\]
\[
\frac{dD_{3j}(t, t + \tau)}{d\tau} = -\Gamma(2N(r) + 1) D_{3j}(t, t + \tau) - \\
\frac{-\Omega_R}{2}(e^{i\omega_f(t+\tau)} + e^{-i\omega_f(t+\tau)})(D_{2j}(t, t + \tau)e^{i\omega_f(t+\tau)} + D_{1j}(t, t + \tau)e^{-i\omega_f(t+\tau)}),
\]

with the only difference that the inhomogeneity \(-\Gamma/2\) disappears due to the subtraction of the mean. For the resonant case \(\omega_f = \omega_0\) these correlation functions can be decomposed in the same way as the components \(X_i(t)\) before, i.e. \(D_{ij}(t, t + \tau) = \sum_{l=-\infty}^{+\infty} D_{ij}^{(l)}(t, \tau) e^{il\omega_f(t+\tau)}, \quad i, j = 1, 2, 3,\)

and the Laplace transforms \(\tilde{D}_{ij}^{(l)}(t, z) = \int_{0}^{\infty} e^{-z\tau} D_{ij}^{(l)}(t, t + \tau) d\tau\) of the components \(D_{ij}^{(l)}(t, \tau)\) will satisfy the following set of equations

\[
\begin{align*}
z\tilde{D}_{1j}^{(l)}(t, z) + \left( \frac{\Gamma}{2}(2N(r) + 1) + il\omega_f \right) \tilde{D}_{1j}^{(l)}(t, z) - \Gamma M(r, \theta) \tilde{D}_{2j}^{(l)}(t, z) + \\
\frac{\delta_\alpha}{2} \left( \tilde{D}_{1j}^{(l-1)}(t, z) + \tilde{D}_{1j}^{(l+1)}(t, z) \right) - \Omega_R \left( \tilde{D}_{3j}^{(l-2)}(t, z) + \tilde{D}_{3j}^{(l)}(t, z) \right) = C_{1j}, \quad \tag{21}
\end{align*}
\]
\[
\begin{align*}
z\tilde{D}_{2j}^{(l)}(t, z) + \left( \frac{\Gamma}{2}(2N(r) + 1) + il\omega_f \right) \tilde{D}_{2j}^{(l)}(t, z) - \Gamma M^*(r, \theta) \tilde{D}_{1j}^{(l)}(t, z) - \\
- \frac{\delta_\alpha}{2} \left( \tilde{D}_{2j}^{(l-1)}(t, z) + \tilde{D}_{2j}^{(l+1)}(t, z) \right) - \Omega_R \left( \tilde{D}_{3j}^{(l+2)}(t, z) + \tilde{D}_{3j}^{(l)}(t, z) \right) = C_{2j}, \quad \tag{22}
\end{align*}
\]
\[
\begin{align*}
z\tilde{D}_{3j}^{(l)}(t, z) + \left( \Gamma(2N(r) + 1) + il\omega_f \right) \tilde{D}_{3j}^{(l)}(t, z) + \\
+ \frac{-\Omega_R}{2} \left( \tilde{D}_{2j}^{(l-2)}(t, z) + \tilde{D}_{2j}^{(l)}(t, z) + \tilde{D}_{1j}^{(l)}(t, z) + \tilde{D}_{1j}^{(l+2)}(t, z) \right) = C_{3j}, \quad \tag{23}
\end{align*}
\]

where

\[
C_{ij} = \begin{bmatrix}
0 & -2X_3^{(l)}(t) & X_1^{(l)}(t) \\
2X_3^{(l)}(t) & 0 & -X_2^{(l)}(t) \\
-X_1^{(l)}(t) & X_2^{(l)}(t) & 0
\end{bmatrix}.
\]
In the steady state \((t \to \infty)\) only a handful of components \(\bar{D}_{ij}^{(l)}(t,z)\) contribute to \(W_\alpha(\omega)\), and the absorption spectrum \((14)\) can be expressed as

\[
W_\alpha(\omega) = C_0 \omega \text{Re} \lim_{t \to \infty} [R^2 \bar{D}_{33}^{(0)}(t,z) + iR (\bar{D}_{13}^{(1)}(t,z) - \bar{D}_{23}^{(-1)}(t,z))]_{z=-i\omega} + \\
+ [ (\bar{D}_{21}^{(0)}(t,z) - \bar{D}_{11}^{(2)}(t,z)) + iR \bar{D}_{31}^{(1)}(t,z) ]_{z=-i(\omega + \omega_f)} + \\
+ [ (\bar{D}_{12}^{(0)}(t,z) - \bar{D}_{22}^{(-2)}(t,z)) - iR \bar{D}_{32}^{(-1)}(t,z) ]_{z=-i(\omega - \omega_f)},
\]

(25)

where \(C_0 = \frac{\hbar}{2} \Omega_R^{(p)}\) and \(R = \delta_\alpha^{(p)} / \Omega_R^{(p)}\).

5. Numerical Results

Equations \((9)-(11)\) and \((21)-(23)\) were solved numerically in the steady state limit \((t \to \infty)\) for the case of the laser field frequency \(\omega_f\) being in resonance with the atomic transition frequency \(\omega_0\). The only approximation made in the course of these calculations consisted in the truncation of the number of harmonic amplitudes \(X_{ij}^{(l)}(t)\) and \(D_{ij}^{(l)}(t,z)\) involved, precisely as was done before in, for example, [7]. It was already found [1] in the case of a two-level system with broken symmetry interacting with non-squeezed vacuum reservoir \(r\) that for \(\delta_\alpha \neq 0\) a peculiar absorption-amplification structure arises in the absorption spectrum around the frequency \(\omega = \Omega_R\), as shown in Fig.1. The left-hand broadband part of this structure for \(\omega < \Omega_R\) is negative, which feature can be interpreted, as usual, as the amplification of the probe field. The amplitude of the amplification-absorption structure grows steadily with the increase of the asymmetry violation parameter \(\delta_\alpha\), while the position of the structure is defined mostly by the Rabi frequency \(\Omega_R\) and also depends weakly on the symmetry violation parameter \(\delta_\alpha\), in the sense that it drifts slowly leftward as this parameter increases. It was found that the increase in the degree of the vacuum reservoir squeezing \(r\) results in the decrease of the absorption-amplification spectrum amplitude, see Fig.2, down to the point of its total disappearance. The shape of the absorption-amplification spectrum changes dramatically along with the variation of the squeezed field reference phase \(\theta\). For positive \(\theta\) the spectrum transforms quickly into purely amplifying spectrum centered at \(\omega \approx \Omega_R\) and back to the absorption-amplification spectrum shown in Fig.1 as the reference phase increases from 0 to \(\pi\), while for negative \(\theta\) the spectrum transforms into purely absorptive spectrum centered at \(\omega \approx \Omega_R\) and goes back to the absorption-amplification spectrum shown in Fig.1 as the reference phase progresses from 0 to \(-\pi\), see Fig.3. The amplitude of the spectrum at \(\omega = \Omega_R\) varies with the reference phase \(\theta\) as depicted in Fig.4. It is seen that for small \(r\) both spectra, the amplifying and the absorptive one, attain their respective extrema at \(\omega = \Omega_R\) for \(\theta = \pi/2\) and \(\theta = -\pi/2\) correspondingly. For larger \(r\) the positions of both extrema shift toward \(\theta = \pm \pi\). For any fixed \(\theta \neq 0\) the amplitude of the spectrum first increases with the increase of \(r\). Then, upon further increase of \(r\), the amplitude of the spectral peak decreases down to its complete disappearance for large enough \(r\).
In conclusion, we obtained and analyzed the weak probe EM field absorption-amplification spectrum for a simple two-level quantum system with broken inversion symmetry possessing non-equal permanent diagonal dipole moment matrix elements, driven by semiclassical monochromatic high-frequency EM field at its resonant frequency $\omega_0$ and damped by broadband squeezed vacuum reservoir. It was shown that, depending on the reference phase of the squeezed field $\theta$, this system can either absorb or amplify low-frequency weak probe field around Rabi frequency $\Omega_R$ of the driving field. At the same time, the increase in the degree of squeezing $r$
leads eventually to the decrease of the spectral amplitude and disappearance of the absorption-amplification spectrum around \( \omega = \Omega_R \). Nevertheless, for reasonably moderate squeezing this phenomenon may provide one with potentially useful practical ability to switch between absorption and amplification of the low-frequency EM radiation, especially in the terahertz range, by altering the reference phase of the squeezed field. A preferable plausible way toward practical implementation of this method for the terahertz radiation manipulation may consist in the building of the aforesaid quantum two-level system with broken inversion symmetry on the basis of semiconductor quantum dots designed and manufactured with desired externally controllable symmetry properties. It is also worth noting that the results of this study can be generalized for the case of broadband squeezed thermal reservoir characterized by the density operator

\[
\rho^{vac,sq}_B = \hat{S}(\xi)\rho_{th} B \hat{S}^+(\xi), \quad \rho_{th} B = \frac{e^{-\beta H_0(B)}}{S_{B\rho_{th}}(\beta)},
\]

(see [6]). Such a generalization only requires substitution

\[
N(r) \rightarrow \bar{N}(r,\omega_f), \quad M(r,\theta) \rightarrow \bar{M}(r,\theta,\omega_f), \quad \bar{N}(r,\omega_f) = \frac{1}{e^{\beta \omega_f} - 1}, \quad \bar{M}(r,\theta,\omega_f) = (2\bar{n}(\omega_f) + 1) \cosh(r) \sinh(r) \exp(i\theta),
\]

throughout all of the equations. But, in reality, \( \bar{n}(\omega_f) \approx 0 \) for the typical range of the laser field frequencies \( \omega_f \) under normal laboratory conditions. Therefore, this substitution can not alter the discussed phenomena significantly.

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