The steady state population inversion of multiple Ξ-type atoms by the squeezed vacuum in a waveguide

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We study the dynamics of multiple Ξ-type atoms driven by a squeezed vacuum reservoir in a quasi-one-dimensional waveguide. We show that in the steady state the atoms can be in a pure state, and the complete population inversion can occur when the direction of dipole moment is properly set. We also study the case when the atoms can couple to each other by the dipole-dipole interaction and we show that the steady state is the direct product of steady state of single-atom case. The result here may be used to study the two-photon laser or collective atomic effect.

I. INTRODUCTION

The concept of population inversion is of fundamental importance in laser physics because the production of population inversion is a key step in the workings of a standard laser. However, a population inversion can never exist for a system at thermal equilibrium because of the spontaneous emission. The achievement of population inversion therefore requires pushing the system into a non-equilibrated state [1]. Thus, the spontaneous emission must be inhibited in order to maintain the population inversion in a steady state. In 1946, Purcell showed that the spontaneous decay rate of an emitter can be modified by engineering the electromagnetic bath environment with which the emitters interact [2]. One famous example of bath engineering is the squeezed vacuum which leads to many novel effects and techniques in quantum optics and atomic spectroscopy. The reduction of quantum fluctuations below vacuum level by the squeezed vacuum yields many interesting phenomenons, for example, the suppression of dephasing rate in one direction and enhancment in the other for a two-level emitter [3–10], and the subnatural linewidth of resonance fluorescence [11, 12]. The entanglement nature of the squeezed vacuum also leads to interesting results like pairwise excitation of atomic states [13–15]. In 1993, Ficek et. al. studied the dynamical properties of a three-level atom in the squeezed vacuum where they showed that a three-level atom in the cascade configuration coupled to squeezed modes can reach steady state level populations inversion relative to ordinary laser spectroscopy [16, 17].

Although the squeezing of all modes in three dimensions is difficult, the squeezing of modes in one dimension is experimentally feasible. Suppression of the radiative decay of atomic coherence and the linewidth of the resonance fluorescence have been experimentally demonstrated in a 1D microwave transmission line coupled to single artificial atom [12, 25–29]. Recently, photon transport in a one-dimensional (1D) waveguide coupled to quantum emitters (well known as "waveguide-QED")

has attracted much attention due to its possible applications in quantum device and quantum information [18–24]. Here, we consider multiple Ξ-type atoms coupled to the broadband squeezed vacuum in the quasi-one-dimensional waveguide where we assume that all modes in one dimension are squeezed. We show that for a single atom, the steady state is the superposition of the ground state and the second excited state and the population inversion can approach 100%. This result is similar to the result in the cavity case shown in Ref. [16]. Furthermore we also mathematically prove that this result can be generalized to arbitrary amount of atoms coupled to each other through dipole-dipole interaction, which may be a scenario for studying the subradiance effect.

The paper is organized as follows. In Sec. II, we derive the master equation of three-level atoms coupled to squeezed vacuum in 1D waveguide. In Sec. III, we derive the steady state of a single three-level atom driven by squeezed vacuum in 1D waveguide. In Sec. IV, we consider the case with multiple atoms which can have dipole-dipole interactions. Finally, we summarize the result.

In contrast to the 3D case, squeezing in 1D is more experimentally feasible. Thus, Ficek's proposl on the steady state popoulation inversion seems to be experimentally implementable. However, there are still three issues in Ficek's proposal: first, the steady state population distribution cannot be controlled because the atomic decay rate is not tunable; second, only a single independent atom is considered to be interacting with the squeezed vacuum, while the dipole-dipole interaction between atoms may affect the steady state population distribution; third, the effect of squeezing source is not included as discussed in Ref.[15]. Our goal is to fix these issues.

In our study, we consider multiple Ξ-type atoms coupled to the broadband squeezed vacuum in the quasi-one-dimensional waveguide where all resonant modes can be technically squeezed. We show that for a single atom, it can always reach a population inversion of almost 100% or any other ratio as long as its dipole direction is prop-

erly set. We also mathematically prove that this result can be generalized to arbitrary amount of atoms coupled to each other through dipole-dipole interaction, which may be a scenario for studying the subradiance effect.

II. MASTER EQUATION OF THREE-LEVEL ATOMS IN THE SQUEEZED VACUUM

In this section, we consider a scenario where N_a Ξ-type atoms are located inside a perfect rectangular waveguide with the squeezed vacuum injected from both ends, as shown in Fig. 1(a). The atomic electronic structure is shown in Fig. 1(b) where the atomic states are labeled as $|a\rangle$, $|b\rangle$, $|c\rangle$ from the excited state to the ground state. Different from the free space, the square waveguide can only support certain photon modes, i.e., TE_{mn} and TM_{mn} modes with cutoff frequency $c\sqrt{(\frac{m\pi}{a})^2 + (\frac{n\pi}{b})^2}$, where $a \times b$ is the dimension of the waveguide's cross section. For the rectangular waveguide, there is no TM_{01} or TM_{10} mode. Assuming $b < a$, TE_{10} is the ground mode with the lowest cutoff frequency. Thus, the atom is only coupled to TE_{10} mode as long as $\frac{\pi c}{a} < \omega_{ab}, \omega_{bc} < \frac{\pi c}{b}, \frac{2\pi c}{a}$. We assume that $\omega_{ac} = 2\omega_0$ where ω_0 is the center frequency of the broadband squeezed vacuum. ω_{ab} and ω_{bc} are not equal but they are still within the bandwidth of the squeezed vacuum. We assume that the squeezed vacuum bandwidth is much larger than $|\omega_{ab} - \omega_{bc}|$ so it forms a squeezed vacuum reservoir.

The atom-field system is described by the Hamiltonian

$$H = H_A + H_F + H_{AF} \qquad (1)$$

where $H_A = \sum_{l=1}^{N_a} \sum_{e=a,b,c} \hbar\omega_{e,l} |e_l\rangle \langle e_l|$ is the atomic Hamiltonian, and $|e_l\rangle$ is the energy state of the lth atom with energy $\hbar\omega_{e,l}$. The Hamiltonian of the EM field is $H_F = \sum_{ks} \hbar\omega_{ks}(\hat{a}_{ks}^\dagger \hat{a}_{ks} + \frac{1}{2})$ where $\hat{a}_{ks}$ and $\hat{a}_{ks}^\dagger$ are the annihilation and creation operators of the filed mode with wavevector k, polarization s (in waveguide, it represents TE_{mn} or TM_{mn}), and frequency $\omega_{k,s}$. The interaction Hamiltonian in electric-dipole approximation is $H_{AF} = -i\hbar \sum_{ks} \sum_{i=1,2} \sum_{l=1}^{N_a} [\boldsymbol{\mu}_{l,i} \cdot \boldsymbol{u}_{ks}(\boldsymbol{r}_{l,i}) S_{l,i}^+ \hat{a}_{ks} + \boldsymbol{\mu}_{l,i}^* \cdot \boldsymbol{u}_{ks}(\boldsymbol{r}_{l,i}) S_{l,i}^- \hat{a}_{ks} - H.c.]$ where $\boldsymbol{\mu}_{l,i}$ is the electric dipole moment for ith transition of the lth atom, where $i = 1$ denotes the transition from $|a\rangle$ to $|b\rangle$, and $i = 2$ denotes the transition from $|b\rangle$ to $|c\rangle$. $S_{l,i}^+$ and $S_{l,i}^-$ are the raising and lowering operator for transition i of the lth atom. The mode function of the squeezed vacuum in TE_{10} mode is given by

$$\boldsymbol{u}_{ks}(\boldsymbol{r}_i) = \sqrt{\frac{\omega_{ks}}{2\epsilon_0 \hbar V}} \boldsymbol{x} e^{i\boldsymbol{k}\cdot(\boldsymbol{r}_i - \boldsymbol{o}_{ks})} \qquad (2)$$

where $\boldsymbol{o}_{ks}$ is a phenomenological parameter which includes the effects of the initial phase and the position of the squeezing source[15]. The correlation functions for the squeezed vacuum are[30]:

$$
\begin{aligned}
\left\langle a_{k,s}^\dagger a_{k',s'} \right\rangle &= \sinh^2 r \delta_{k'k}\delta_{ss'} \\
\left\langle a_{k,s} a_{k',s'}^\dagger \right\rangle &= \cosh^2 r \delta_{k'k}\delta_{ss'} \\
\left\langle a_{k,s}^\dagger a_{k',s'}^\dagger \right\rangle &= -e^{-i\theta}\cosh(r)\sinh(r)\delta_{k',2k_0-k}\delta_{ss'} \\
\left\langle a_{k,s} a_{k',s'} \right\rangle &= -e^{i\theta}\cosh(r)\sinh(r)\delta_{k',2k_0-k}\delta_{ss'}
\end{aligned}
\qquad (3)
$$

For simplicity, we can set the squeezing parameter $\theta = 0$, and the dipole moments of all atoms to align along the same direction. The dynamics of the atomic system can be described by the following master equation (See Appendix A for details of derivation):

$$
\begin{aligned}
\frac{d\rho^S}{dt} = &-i\sum_{ijkl}\Lambda_{ijkl}[S_{i,j}^+ S_{k,l}^-, \rho^S]e^{i(\omega_j - \omega_l)t} - \frac{1}{2}\sum_{ijkl}\gamma_{ijkl}(1+N)(\rho^S S_{i,j}^+ S_{k.l}^- + S_{i,j}^+ S_{k.l}^- \rho^S - 2S_{k.l}^- \rho^S S_{i,j}^+)e^{i(\omega_j - \omega_l)t} \\
&- \frac{1}{2}\sum_{ijkl}\gamma_{ijkl}N(\rho^S S_{i,j}^- S_{k.l}^+ + S_{i,j}^- S_{k,l}^+ \rho^S - 2S_{k,l}^+ \rho^S S_{i,j}^-)e^{-i(\omega_j - \omega_l)t} \\
&- \frac{1}{2}\sum_{\alpha=\pm}\sum_{ijkl}\gamma'_{ijkl}M e^{2\alpha i k_{0z}R}e^{i\alpha(\omega_j + \omega_l - 2\omega_0)t}(\rho^S S_{i,j}^\alpha S_{k.l}^\alpha + S_{i,j}^\alpha S_{k,l}^\alpha \rho^S - 2S_{k,l}^\alpha \rho^S S_{i,j}^\alpha)
\end{aligned}
\qquad (4)
$$

where $N = \sinh(r)^2$, $M = \sinh(r)\cosh(r)$, and the coefficients are

$$
\begin{aligned}
\gamma_{ijkl} &= \sqrt{\gamma_j\gamma_l}\cos(k_{0z}r_{ik}) \\
\Lambda_{ijkl} &= \frac{\sqrt{\gamma_j\gamma_l}}{2}\sin(k_{0z}r_{ik}) \\
\gamma'_{ijkl} &= \sqrt{\gamma_j\gamma_l}\cos[k_{0z}(r_i + r_k)]
\end{aligned}
\qquad (5)
$$

where subscript i, k label the atom index, $j(l)$ labels the transitions of the ith(kth) atom. γ_j is the decay rate for transition j in ordinary vacuum in the waveguide with the dipole moment pointed in the same direction. Contrary to the free space where the decay rate is independent of the dipole direction, the anisotropy character of the waveguide results in the sensitivity of atomic decay rate
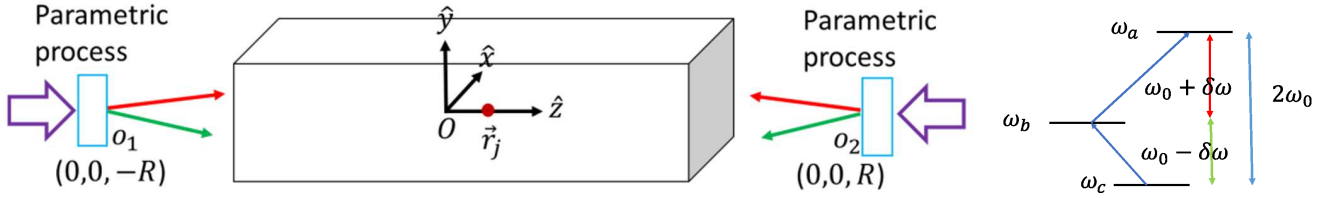
Fig. 1: (a) Schematic setup: A Ξ-type atom is located inside the waveguide with the broadband squeezed vacuum incident from both ends. (b)The energy structure of the three level atom. Transition $|a\rangle \to |c\rangle$ is forbidden and $\omega_{ac} = 2\omega_0$ where ω_0 is the center frequency of the squeezed vacuum. ω_{ab} and ω_{bc} differ by a small amount $2\delta\omega_0$ and they are within the bandwidth of the squeezed vacuum reservoir.

to its dipole direction. Thus, $\gamma_j = \gamma_{jmax}\cos\theta_j$ where γ_{jmax} is the maximum value of γ_j when the dipole is along the x axis, and $\cos\theta_j$ is the angle between the x axis and the dipole j's direction.

III. STEADY STATE OF A SINGLE ATOM

In this section, we will study the steady state of a single atom in the squeezed vacuum reservoir. For a single three

level atom, we have $r_i = r_k$, and for simplicity we set $R = r_i = 0$. Then the steady state of Eq.(4) can be derived by re-writing it as:

$$\dot{\rho}_{aa} = -\gamma_1 ch^2 \rho_{aa} + \gamma_1 sh^2 \rho_{bb} - \frac{1}{2}\sqrt{\gamma_1\gamma_2}chsh(e^{-i(\omega_1+\omega_2-2\omega_0)t}\rho_{ac} + e^{i(\omega_1+\omega_2-2\omega_0)t}\rho_{ca}) \tag{6a}$$

$$\dot{\rho}_{bb} = \gamma_1(ch^2\rho_{aa} - sh^2\rho_{bb}) + \gamma_2(sh^2\rho_{cc} - ch^2\rho_{bb}) + \sqrt{\gamma_1\gamma_2}chsh(e^{-i(\omega_1+\omega_2-2\omega_0)t}\rho_{ac} + e^{i(\omega_1+\omega_2-2\omega_0)t}\rho_{ca}) \tag{6b}$$

$$\dot{\rho}_{cc} = \gamma_2 ch^2\rho_{bb} - \gamma_2 sh^2\rho_{cc} - \frac{1}{2}\sqrt{\gamma_1\gamma_2}chsh(e^{-i(\omega_1+\omega_2-2\omega_0)t}\rho_{ac} + e^{i(\omega_1+\omega_2-2\omega_0)t}\rho_{ca}) \tag{6c}$$

$$e^{-i(\omega_1+\omega_2-2\omega_0)t}\dot{\rho}_{ac} + e^{i(\omega_1+\omega_2-2\omega_0)t}\dot{\rho}_{ca} = -\frac{1}{2}(\gamma_1 ch^2 + \gamma_2 sh^2)(e^{-i(\omega_1+\omega_2-2\omega_0)t}\rho_{ac} + e^{i(\omega_1+\omega_2-2\omega_0)t}\rho_{ca})$$
$$- \sqrt{\gamma_1\gamma_2}shch(\rho_{aa} - 2\rho_{bb} + \rho_{cc}) \tag{6d}$$

$$e^{i(\omega_0-\omega_1)t}\dot{\rho}_{ab} + e^{-i(\omega_0-\omega_1)t}\dot{\rho}_{ba} = -\frac{1}{2}((\gamma_1+\gamma_2)ch^2 + \gamma_1 sh^2 - \gamma_1 chsh)(e^{i(\omega_0-\omega_1)t}\rho_{ab} + e^{-i(\omega_0-\omega_1)t}\rho_{ba})$$
$$- \frac{1}{2}\sqrt{\gamma_1\gamma_2}(ch - 2sh)sh(e^{-i(\omega_0-\omega_2)t}\rho_{cb} + e^{i(\omega_0-\omega_2)t}\rho_{bc}) \tag{6e}$$

$$e^{-i(\omega_0-\omega_2)t}\dot{\rho}_{cb} + e^{i(\omega_0-\omega_2)t}\dot{\rho}_{bc} = \frac{1}{2}\sqrt{\gamma_1\gamma_2}(2ch - sh)ch(e^{i(\omega_0-\omega_1)t}\rho_{ab} + e^{-i(\omega_0-\omega_1)t}\rho_{bc})$$
$$- \frac{1}{2}((\gamma_1+\gamma_2)sh^2 + \gamma_2 ch^2 - 2\gamma_2 chsh)(e^{-i(\omega_0-\omega_2)t}\rho_{cb} + e^{i(\omega_0-\omega_2)t}\rho_{bc}) \tag{6f}$$

where $ch = \cosh(r)$, $sh = \sinh(r)$, and $\gamma_1 = \gamma_{ab}(\gamma_2 = \gamma_{bc})$ is the decay rate from $|a\rangle$ to $|b\rangle (|b\rangle$ to $|c\rangle)$ in ordinary vacuum due to the waveguide modes. Equations (6e)-(6f) are for off diagonal elements ρ_{ab}, ρ_{bc}. The steady state of these two equations is $\rho_{ab} = \rho_{bc} = 0$ since they are homogeneous linear equations. The first four equations Eq.(6a)-(6d) also have a steady state solution when they are combined with the constraints $\rho_{aa} + \rho_{bb} + \rho_{cc} = 1$ and $\omega_1 + \omega_2 = 2\omega_0$. It is also worth noting that Eq.(6a)-(6d) are independent of $\delta\omega$, so the difference between ω_{ab} and ω_{bc} doesn't influence the steady state for the single

atom case if both ω_{ab} and ω_{bc} are within the squeezing bandwidth. When $\gamma_1 \neq \gamma_2$, the steady state solution is:

$$|\psi_{ss}\rangle = \frac{sh\sqrt{\gamma_2}}{\sqrt{ch^2\gamma_1 + sh^2\gamma_2}}|a\rangle - \frac{ch\sqrt{\gamma_1}}{\sqrt{ch^2\gamma_1 + sh^2\gamma_2}}|c\rangle \tag{7}$$

which is a superposition state of $|a\rangle$ and $|c\rangle$. Since there is no population in the state $|b\rangle$, population inversion can always occur between states $|a\rangle$ and $|b\rangle$ in the steady state. If $\tanh r > \sqrt{\frac{\gamma_1}{\gamma_2}}$ population inversion can also occur between the state $|a\rangle$ and $|c\rangle$. This result is similar to

the result in the cavity case [16]. However, different from the result in the cavity case, the population inversion in the 1D waveguide case can approach 100% if $\gamma_2 \gg \gamma_1$. As we mentioned before, γ_1 and γ_2 are controllable if we can control the alignment of the atomic dipole moments, so the resultant population inversion can also be controlled.

The steady state population distribution for different ratios of $\frac{\gamma_{ab}}{\gamma_{bc}}$ is shown in Fig. 2(a). The mechanism of this population inversion can be interpreted with the help of Fig. 2(c). Figure 2(c) shows that the direct transition between $|a\rangle$, $|b\rangle$, and $|c\rangle$ are allowed just like the thermal reservoir case. However, in the squeezed vacuum, there is additional paths for population flow: atom in any of these three states can evolve into the other two through an intermediate "state" ρ_{ac}. Although ρ_{ac} is actually an off-diagonal element rather than a state, it can be used to elucidate our idea. When $\gamma_{ab} \ll \gamma_{bc}$, the transition $|a\rangle \to |b\rangle$ is negligible compared with γ_{bc} and $\sqrt{\gamma_{ab}\gamma_{bc}}$. Thus the atom in the state $|c\rangle$ can be excited to $|a\rangle$ through $|c\rangle \to |b\rangle \to \rho_{ac} \to |a\rangle$, but $|a\rangle$ can not decay back to $|c\rangle$, which results in the population trapping in the $|a\rangle$. This phenomenon is similar to the coherent trapping, but here we achieve the trapping for Ξ structure with the squeezed vacuum reservoir, which cannot be realized with coherent pump due to spontaneous emission. What's more, since only a few modes are allowed in the waveguide with particular polarizations [15], $\gamma_{ab} \ll \gamma_{bc}$ can be achieved if we can properly manipulate the orientation of the atom. Since it is hard to achieve perfect squeezing with $M = \sqrt{N(N+1)}$ in experiments, we also study the effect of different values of M on the steady state population with parameters $\gamma_{ab} = \frac{1}{4}\gamma_{bc}$ and $r = 1$, which is shown in Fig. 2(b). In general, all three energy levels have population. Although the steady state population distribution is very sensitive to the value of M, the population inversion between $|a\rangle$ and $|b\rangle$ still holds for $M = 0.8\sqrt{N(N+1)}$. Only when M is larger than 0.95 can the population inversion occurs between the state $|a\rangle$ and the state $|c\rangle$.

IV. STEADY STATE OF MULTIPLE ATOMS

In the last section, we demonstrate that a single atom can be driven to the steady state population inversion in the squeezed vacuum reservoir. However, with Eq.(5), this result can not be generalized to the multi-atom case since $\gamma'_{ijij} = \sqrt{\gamma_j\gamma_j}\cos[2k_{0z}r_i]$. The squeezing term in Eq.(4) vanishes for atoms located around $r_i = \frac{\pi}{4k_{0z}} + \frac{n\pi}{2k_{0z}}$. Thus, for a group of randomly located atoms, if we want to achieve steady state population inversion in the squeezed vacuum, we need to modify our scheme. Here
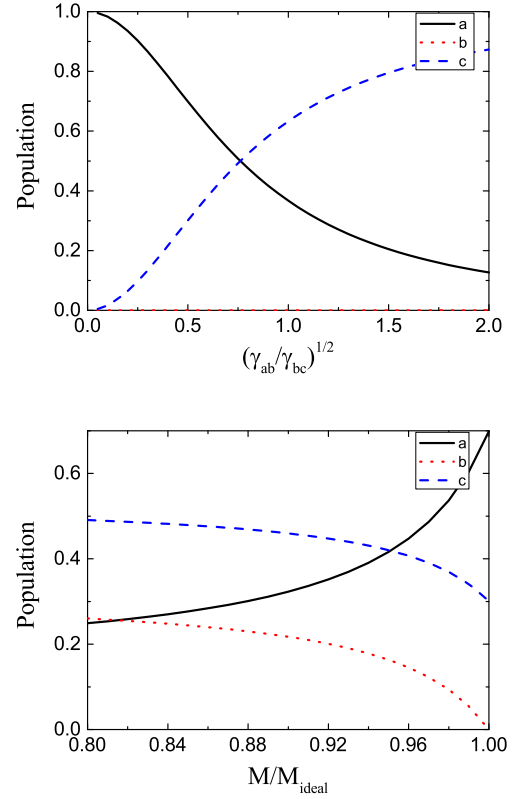

Fig. 2: (a) The steady state population distribution for different μ_{ab} and μ_{bc}. The squeezing parameter $r = 1$. (b) The steady state population distribution for non-ideal squeezed vacuum which is characterized by the ratio of M and $\sqrt{N(N+1)}$.
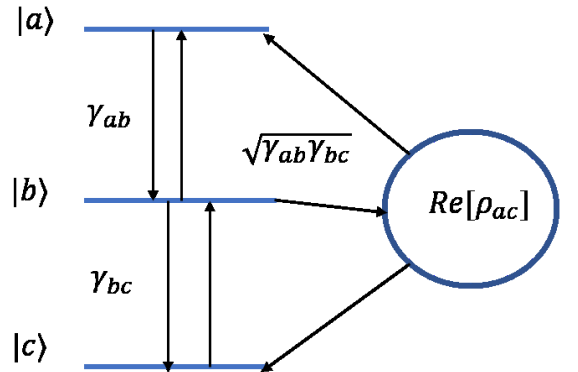

Fig. 3: The allowed population flow in the squeezed vacuum. The squeezing parameter $r = 1$, and $\gamma_{ab} = \frac{1}{4}\gamma_{bc}$.

we consider the following correlation functions:

$$\left\langle a_{\boldsymbol{k},s}^{\dagger} a_{\boldsymbol{k}',s'} \right\rangle = \sinh^2 r \delta_{\boldsymbol{k}'\boldsymbol{k}} \delta_{ss'}$$

$$\left\langle a_{\boldsymbol{k},s} a_{\boldsymbol{k}',s'}^{\dagger} \right\rangle = \cosh^2 r \delta_{\boldsymbol{k}'\boldsymbol{k}} \delta_{ss'}$$

$$\left\langle a_{\boldsymbol{k},s}^{\dagger} a_{\boldsymbol{k}',s'}^{\dagger} \right\rangle = -e^{-i\theta} \cosh(r) \sinh(r) \delta_{\boldsymbol{k}',-(2\boldsymbol{k}_0-\boldsymbol{k})} \delta_{ss'}$$

$$\langle a_{\boldsymbol{k},s} a_{\boldsymbol{k}',s'} \rangle = -e^{i\theta} \cosh(r) \sinh(r) \delta_{\boldsymbol{k}',-(2\boldsymbol{k}_0-\boldsymbol{k})} \delta_{ss'}$$

$$(8)$$

which indicates that photons are entangled with those from the opposite direction, then the coefficients in the the master equation Eq.(4) becomes

$$\gamma_{ijkl} = \sqrt{\gamma_i \gamma_k} \cos(k_{0z} r_{jl})$$

$$\Lambda_{ijkl} = \frac{\sqrt{\gamma_i \gamma_k}}{2} \sin(k_{0z} r_{jl}) \qquad (9)$$

$$\gamma'_{ijkl} = \sqrt{\gamma_i \gamma_k} \cos[k_{0z}(r_{jl})]$$

which is the traditionally studied master equation for atoms in squeezed reservoir [13]. With the phase of the squeezing source included, the detailed derivation of the coefficients shown in Eq. (9) is given in Appendix B. Based on the master equation in Eq. (4) with coefficients given by Eq. (9), we can show that a single atom can reach population inversion anywhere in the waveguide, as discussed in the Sec. III. When there are multiple atoms in the waveguide where the dipole-dipole interaction should be considered, our calculation shows that the population inversion can still occur for all the atoms. In fact, the final state is just the direct product of the steady state of independent atoms, and this will be proved with the mathematical induction. Firstly, we need to apply the rotating wave approximation on Eq.(4):

$$\begin{aligned}
\frac{d\rho^S}{dt} = & -i\sum_{i,k,j} \Lambda_{ijkl}[S_{i,j}^+ S_{k,j}^-, \rho^S] - \frac{1}{2}\sum_{i,j,k} \gamma_{ijkl}(1+N)(\rho^S S_{i,j}^+ S_{k,j}^- + S_{i,j}^+ S_{k,j}^- \rho^S - 2S_{k,j}^- \rho^S S_{i,j}^+) \\
& -\frac{1}{2}\sum_{i,j,k} \gamma_{ijkl} N(\rho^S S_{i,j}^- S_{k,j}^+ + S_{i,j}^- S_{k,j}^+ \rho^S - 2S_{k,j}^+ \rho^S S_{i,j}^-) \\
& -\frac{1}{2}\sum_{\alpha=\pm}\sum_{i,k,j\neq l} \gamma'_{ijkl} M(\rho^S S_{i,j}^\alpha S_{k,l}^\alpha + S_{i,j}^\alpha S_{k,l}^\alpha \rho^S - 2S_{k,l}^\alpha \rho^S S_{i,j}^\alpha)
\end{aligned} \qquad (10)$$

Assume that the steady state of N-atom system is $\rho^S = \rho_1 \rho_2 ... \rho_N$ where $\rho_i = (A|a_i\rangle + C|c_i\rangle)(A\langle a_i| + C\langle c_i|)$ and $A = \frac{sh\sqrt{\gamma_2}}{\sqrt{ch^2\gamma_1 + sh^2\gamma_2}}, C = -\frac{ch\sqrt{\gamma_1}}{\sqrt{ch^2\gamma_1 + sh^2\gamma_2}}$. Then for $N+1$-atom case, the extra terms induced by the $(N+1)th$ atom on the right hand side of Eq.(10) are composed of three parts: $i = k = N+1$ terms, $i = N+1, k = $

$1, 2, \cdots, N$ terms, and $i = 1, 2, \cdots, N, k = N+1$ terms. The $i = k = N+1$ terms are the exact terms for the single independent atom of the $(N+1)th$ atom, so the result is 0. The terms with $i = N+1, k = 1, 2, \cdots, N$ are

$$\begin{aligned}
& -i\sum_{j,k} \Lambda_{N+1,j,k,j}[S_{N+1,j}^+ S_{k,j}^-, \rho^S] - \frac{1}{2}\sum_{j,k} \gamma_{N+1,j,k,j}(ch^2)(\rho^S S_{N+1,j}^+ S_{k,j}^- + S_{N+1,j}^+ S_{k,j}^- \rho^S - 2S_{k,j}^- \rho^S S_{N+1,j}^+) \\
& -\frac{1}{2}\sum_{j,k} \gamma_{N+1,j,k,j} sh^2 (\rho^S S_{N+1,j}^- S_{k,j}^+ + S_{N+1,j}^- S_{k,j}^+ \rho^S - 2S_{k,j}^+ \rho^S S_{N+1,j}^-) \\
& -\frac{1}{2}\sum_{\alpha=\pm}\sum_{i,k,j\neq l} \gamma'_{N+1,j,k,j} chsh(\rho^S S_{N+1,j}^\alpha S_{k,l}^\alpha + S_{N+1,j}^\alpha S_{k,l}^\alpha \rho^S - 2S_{k,l}^\alpha \rho^S S_{N+1,j}^\alpha)
\end{aligned} \qquad (11)$$

For the energy shift term (the first term) of Eq.(11), we have

$$S_{N+1,j}^+ S_{k,j}^- \rho^S = \rho_1 ... (S_{k,j}^- \rho_k) ... \rho_N (S_{N+1,j}^+ \rho_{N+1}) = 0$$

$$\rho^S S_{N+1,j}^+ S_{k,l}^- = \rho_1 ... (\rho_k S_{k,l}^-) ... \rho_N (\rho_{N+1} S_{N+1,j}^+) = 0$$
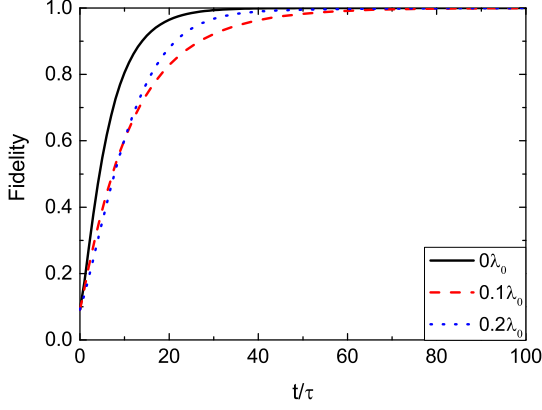
$$(12)$$

Fig. 4: Fidelity evolution for different atomic separation. The atomic separation of 0, $0.1\lambda_0$, $0.2\lambda_0$ are plotted, and τ is the single atom relaxation time in ordinary vacuum.

For the thermal terms (the second and third terms) of Eq.(11), we have:

$$
\begin{aligned}
\rho^S S^+_{N+1,j} S^-_{k,j} &= S^+_{N+1,j} S^-_{k,j} \rho^S = 0 \\
\rho^S S^-_{N+1,j} S^+_{k,j} &= S^-_{N+1,j} S^+_{k,j} \rho^S = 0 \\
S^-_{k,j} \rho^S S^+_{N+1,j} &= \rho_1...(S^-_{k,j}\rho_k)...\rho_N(\rho_{N+1}S^+_{N+1,j}) \\
&= \rho_1...(S^-_{k,1}\rho_k)...\rho_N(\rho_{N+1}S^+_{N+1,1}) \\
&= \rho_1...(A|b_k\rangle)(A\langle a_k| + C\langle c_k|)..\rho_N \\
&\quad \times (A|a_{N+1}\rangle + C|c_{N+1}\rangle)(A\langle b_{N+1}|) \\
S^+_{k,j} \rho^S S^-_{N+1,j} &= \rho_1...(S^+_{k,j}\rho_k)...\rho_N(\rho_{N+1}S^-_{N+1,j}) \\
&= \rho_1...(S^+_{k,2}\rho_k)...\rho_N(\rho_{N+1}S^-_{N+1,2}) \\
&= \rho_1...(C|b_k\rangle)(A\langle a_k| + C\langle c_k|)...\rho_N \\
&\quad \times (A|a_{N+1}\rangle + C|c_{N+1}\rangle)(C\langle b_{N+1}|)
\end{aligned}
\tag{13}
$$

For the squeezed vacuum terms(the forth term), we have:

$$
\begin{aligned}
\rho^S S^\alpha_{N+1,j} S^\alpha_{k,l} &= S^\alpha_{N+1,j} S^\alpha_{k,l}\rho^S = 0, \\
S^+_{k,1} \rho^S S^+_{N+1,2} &= \rho_1...(S^+_{k,1}\rho_k)...\rho_N(\rho_{N+1}S^+_{N+1,2}) = 0, \\
S^+_{k,2} \rho^S S^+_{N+1,1} &= \rho_1...(S^+_{k,2}\rho_k)...\rho_N(\rho_{N+1}S^+_{N+1,1}) \\
&= \rho_1...(C|b_k\rangle)(A\langle a_k| + C\langle c_k|)...\rho_N \\
&\quad \times (A|a_{N+1}\rangle + C|c_{N+1}\rangle)(A\langle b_{N+1}|), \\
S^-_{k,1} \rho^S S^-_{N+1,2} &= \rho_1...(S^-_{k,1}\rho_k)...\rho_N(\rho_{N+1}S^-_{N+1,2}) \\
&= \rho_1...(A|b_k\rangle)(A\langle a_k| + C\langle c_k|)...\rho_N \\
&\quad \times (A|a_{N+1}\rangle + C|c_{N+1}\rangle)(C\langle b_{N+1}|), \\
S^-_{k,2} \rho^S S^-_{N+1,1} &= \rho_1...(S^-_{k,2}\rho_k)...\rho_N(\rho_{N+1}S^-_{N+1,1}) = 0.
\end{aligned}
\tag{14}
$$

Plugging Eq.(12)-(14) into expression (11), we have

$$
\begin{aligned}
&\sum_k \gamma_{N+1,1,k,1}(ch^2)S^+_{k,1}\rho^S S^-_{N+1,1} + \sum_k \gamma_{N+1,2,k,2}sh^2 S^+_{k,2}\rho^S S^-_{N+1,2} + \sum_{\alpha=\pm}\sum_{i,k,j\neq l}\gamma'_{N+1,j,k,l}chsh S^\alpha_{k,l}\rho^S S^\alpha_{N+1,j} \\
&= \sum_k \gamma_{N+1,1,k,1}(ch^2)\rho_1...(A|b_k\rangle)(A\langle a_k| + C\langle c_k|)..\rho_N(A|a_{N+1}\rangle + C|c_{N+1}\rangle)(A\langle b_{N+1}|) \\
&\quad + \sum_k \gamma_{N+1,2,k,2}sh^2 \rho_1...(C|b_k\rangle)(A\langle a_k| + C\langle c_k|)...\rho_N(A|a_{N+1}\rangle + C|c_{N+1}\rangle)(C\langle b_{N+1}|) \\
&\quad + \sum_k \gamma'_{N+1,j,k,l}chsh[\rho_1...(C|b_k\rangle)(A\langle a_k| + C\langle c_k|)...\rho_N(A|a_{N+1}\rangle + C|c_{N+1}\rangle)(A\langle b_{N+1}|) \\
&\quad + \rho_1...(A|b_k\rangle)(A\langle a_k| + C\langle c_k|)...\rho_N(A|a_{N+1}\rangle + C|c_{N+1}\rangle)(C\langle b_{N+1}|)] \\
&= \sum_k(\gamma_{N+1,1,k,1}ch^2 A^2 + \gamma_{N+1,2,k,2}sh^2 C^2 + 2\gamma'_{N+1,j,k,l}chshCA) \\
&\quad \times \rho_1...(|b_k\rangle)(A\langle a_k| + C\langle c_k|)..\rho_N(A|a_{N+1}\rangle + C|c_{N+1}\rangle)(\langle b_{N+1}|).
\end{aligned}
\tag{15}
$$

Since $ch^2 A^2\gamma_{N+1,1,k,1} + sh^2 C^2\gamma_{N+1,2,k,2} + 2chshCA\gamma'_{N+1,j,k,l} = 0$, the extra terms induced by $i = N+1, k = 1 \sim N$ are 0. The same goes for $i = 1 \sim N, k = N+1$. Thus, the direct product state of the steady state of single atom is also the steady state of the multiple atoms. It is interesting that the dipole-dipole interaction between the atoms doesn't have effects on the final steady state. Therefore, just like the single atom case, for multiple atoms, as long as their dipole directions are properly set, a population

inversion of almost 100% can be achieved for each atom.

We also numerically show that steady state of multiple atoms is indeed the direct product of steady state of single atom. Since the cost for numerical simulation increases exponentially as the atom number increases, we only show the fidelity of two-atom state with respect to theoretical steady product state as a function of time in Fig. 4. The system is initially in the ground state. We can see that the system evolves into $(\frac{sh\sqrt{\gamma_2}}{\sqrt{ch^2\gamma_1+sh^2\gamma_2}}|a_1\rangle - \frac{ch\sqrt{\gamma_1}}{\sqrt{ch^2\gamma_1+sh^2\gamma_2}}|c_1\rangle)(\frac{sh\sqrt{\gamma_2}}{\sqrt{ch^2\gamma_1+sh^2\gamma_2}}|a_2\rangle - \frac{ch\sqrt{\gamma_1}}{\sqrt{ch^2\gamma_1+sh^2\gamma_2}}|c_2\rangle)$ no matter what the atomic separation is.

waveguide, with the overall transition frequency $\omega_{ac} = 2\omega_0$. We show that the atoms can evolve into a steady state where each atom is in the superposition state of the ground state and the second excited state even if dipole-dipole interaction between atoms exist. If the decay rate from the second excited state to the first excited state is much smaller than that from the first excited state to the ground state, the population can be almost trapped in the second excited state. This may be useful for two-photon laser.

V. SUMMARY

We study the Ξ-type atoms coupled to a broadband squeezed vacuum reservoir in a quasi-one-dimensional

APPENDIX A: DERIVATION OF EQ.(4)

Here we will show how to derive the master equation Eq. (4). The interaction Hamiltonian is:

$$V(t) = -i\hbar \sum_{\boldsymbol{k}s}[D(t)a_{\boldsymbol{k}s}(t) - D^+(t)a_{\boldsymbol{k}s}^\dagger(t)], \tag{A1}$$

where

$$D(t) = \sum_{l,i}[\boldsymbol{\mu}_{l,i} \cdot \boldsymbol{u}_{\boldsymbol{k},s}(r_{l,i})S_{l,i}^\dagger(t) + \boldsymbol{\mu}_{l,i}^* \cdot \boldsymbol{u}_{\boldsymbol{k},s}(r_{l,i})S_{l,i}^-(t)] \tag{A2}$$

The reduced master equation of atoms in the reservoir is:

$$\begin{aligned}\frac{d\rho^S}{dt} = &-\frac{1}{\hbar^2}\int_0^t d\tau Tr_F\{[V(t),[V(t-\tau),\rho^S(t-\tau)\rho^F\}\\ = &-\frac{1}{\hbar^2}\int_0^t d\tau Tr_F\{V(t)V(t-\tau)\rho^S(t-\tau)\rho^F + \rho^S(t-\tau)\rho^F V(t-\tau)V(t)\\ &-V(t)\rho^S(t-\tau)\rho^F V(t-\tau) - V(t-\tau)\rho^S(t-\tau)\rho^F V(t)\}\end{aligned} \tag{A3}$$

Here we just show how to deal with the first term in Eq.(A3), the remaining terms can be calculated in the same way. For the first term, we have

$$\begin{aligned}&-\frac{1}{\hbar^2}\int_0^t d\tau Tr_F\{V(t)V(t-\tau)\rho^S(t-\tau)\rho^F\}\\ = &\int_0^t d\tau \sum_{\boldsymbol{k}s,\boldsymbol{k}'s'}\{D(t)D(t-\tau)Tr_F[\rho^F a_{ks}(t)a_{k's'}(t-\tau)] - D(t)D^+(t-\tau)Tr_F[\rho^F a_{ks}(t)a_{k's'}^\dagger(t-\tau)]\\ &-D^+(t)D(t-\tau)Tr_F[\rho^F a_{ks}^\dagger(t)a_{k's'}(t-\tau)] + D^+(t)D^+(t-\tau)Tr_F[\rho^F a_{ks}^\dagger(t)a_{k's'}^\dagger(t-\tau)]\}\rho^S(t-\tau)\}.\end{aligned} \tag{A4}$$

Under the rotating wave approximation(RWA), we have

$$-\frac{1}{\hbar^2}\int_0^t d\tau Tr_F\{V(t)V(t-\tau)\rho^S(t-\tau)\rho^F\}$$

$$=\sum_{ijlm}\sum_{\boldsymbol{ks},\boldsymbol{k's'}}\int_0^t d\tau\{\boldsymbol{\mu}_{l,i}\cdot\boldsymbol{u}_{\boldsymbol{ks}}(r_{l,i})S_{l,i}^+e^{i\omega_i t}\boldsymbol{\mu}_{m,j}\cdot\boldsymbol{u}_{\boldsymbol{k's'}}(r_{m,j})S_{m,j}^+e^{i\omega_j(t-\tau)}e^{-i(\omega_{\boldsymbol{ks}}+\omega_{\boldsymbol{k's'}})t+i\omega_{\boldsymbol{k's'}}\tau}[-\sinh(r)\cosh(r)\delta_{\boldsymbol{k'},2\boldsymbol{k}_0-\boldsymbol{k}}\delta_{ss'}]$$

$$-\boldsymbol{\mu}_{l,i}\cdot\boldsymbol{u}_{\boldsymbol{ks}}(r_{l,i})S_{l,i}^+e^{i\omega_i t}\boldsymbol{\mu}_{m,j}^*\cdot\boldsymbol{u}_{\boldsymbol{k's'}}^*(r_{m,j})S_{m,j}^-e^{-i\omega_j(t-\tau)}e^{-i\omega_{\boldsymbol{k's'}}\tau}\cosh^2 r\delta_{\boldsymbol{kk'}}\delta_{ss'}$$

$$-\boldsymbol{\mu}_{l,i}^*\cdot\boldsymbol{u}_{\boldsymbol{ks}}(r_{l,i})S_{l,i}^-e^{-i\omega_i t}\boldsymbol{\mu}_{m,j}\cdot\boldsymbol{u}_{\boldsymbol{k's'}}^*(r_{m,j})S_{m,j}^+e^{i\omega_j(t-\tau)}e^{-i\omega_{\boldsymbol{k's'}}\tau}\cosh^2 r\delta_{\boldsymbol{kk'}}\delta_{ss'}$$

$$-\boldsymbol{\mu}_{l,i}^*\cdot\boldsymbol{u}_{\boldsymbol{ks}}^*(r_{l,i})S_{l,i}^-e^{-i\omega_i t}\boldsymbol{\mu}_{m,j}\cdot\boldsymbol{u}_{\boldsymbol{k's'}}(r_{m,j})S_{m,j}^+e^{i\omega_j(t-\tau)}e^{i\omega_{\boldsymbol{k's'}}\tau}\sinh^2 r\delta_{\boldsymbol{kk'}}\delta_{ss'}$$

$$-\boldsymbol{\mu}_{l,i}\cdot\boldsymbol{u}_{\boldsymbol{ks}}^*(r_{l,i})S_{l,i}^+e^{i\omega_i t}\boldsymbol{\mu}_{m,j}^*\cdot\boldsymbol{u}_{\boldsymbol{k's'}}(r_{m,j})S_{m,j}^-e^{-i\omega_j(t-\tau)}e^{i\omega_{\boldsymbol{k's'}}\tau}\sinh^2 r\delta_{\boldsymbol{kk'}}\delta_{ss'}$$

$$+\boldsymbol{\mu}_{l,i}^*\cdot\boldsymbol{u}_{\boldsymbol{ks}}^*(r_{l,i})S_{l,i}^-e^{-i\omega_i t}\boldsymbol{\mu}_{m,j}^*\cdot\boldsymbol{u}_{\boldsymbol{k's'}}^*(r_{m,j})S_{m,j}^-e^{-i\omega_j(t-\tau)}e^{i(\omega_{\boldsymbol{ks}}+\omega_{\boldsymbol{k's'}})t-i\omega_{\boldsymbol{k's'}}\tau}[-\sinh(r)\cosh(r)\delta_{\boldsymbol{k'},2\boldsymbol{k}_0-\boldsymbol{k}}\delta_{ss'}]\}\rho^S(t-\tau)$$

$$(A5)$$

where l,m are used for labeling different atoms, and i,j are used for transitions within an atom. Here we just calculate the first and second term to show how to get the master equation Eq.(4). Since all atoms are identical, $\omega_{l,i}=\omega_i$, $|\boldsymbol{\mu}_{l,i}|=|\boldsymbol{\mu}_i|$, and $r_{l,i}=r_l$ can be used to simplify Eq.(A5). For simplicity, we define μ_j to be the projection of $\boldsymbol{mu_j}$ on the x axis. For the second term(thermal term), we have

$$-\sum_{k_z}\int_0^t d\tau\boldsymbol{\mu}_{l,i}\cdot\boldsymbol{u}_{\boldsymbol{ks}}(r_l)S_{l,i}^+e^{i\omega_i t}\boldsymbol{\mu}_{m,j}^*\cdot\boldsymbol{u}_{\boldsymbol{k's'}}^*(r_m)S_{m,j}^-e^{-i\omega_j(t-\tau)}e^{-i\omega_{\boldsymbol{k's'}}\tau}\cosh^2 r\rho^S(t-\tau)\delta_{\boldsymbol{kk'}}\delta_{ss'}$$

$$=-\frac{L}{2\pi}e^{i(\omega_i-\omega_j)t}\int_{-\infty}^\infty dk_z\int_0^t d\tau e^{i\omega_j\tau}e^{-i\omega_{k_z}\tau}\frac{\omega_k\mu_i\mu_j}{\epsilon_0 LS\hbar}e^{ik_z(r_l-r_m)}\cosh^2 rS_{l,i}^+S_{m,j}^-\rho^S(t-\tau)$$

$$\approx-\frac{L}{2\pi}e^{i(\omega_i-\omega_j)t}\int_0^\infty dk_z\int_0^t d\tau e^{i\omega_j\tau}e^{-i[\omega_j+c^2 k_{jz}(k_z-k_{jz})/\omega_j]\tau}\frac{\omega_k\mu_i\mu_j}{\epsilon_0 LS\hbar}[e^{ik_z(r_l-r_m)}+e^{-ik_z(r_l-r_m)}]\cosh^2 rS_{l,i}^+S_{m,j}^-\rho^S(t-\tau)$$

$$\approx-\frac{L}{2\pi}e^{i(\omega_i-\omega_j)t}\int_{-k_{0z}}^\infty d\delta k_z\int_0^t d\tau e^{-i\tau c^2 k_{jz}\delta k_z/\omega_j}\frac{\omega_k\mu_i\mu_j}{\epsilon_0 LS\hbar}[e^{i(k_{jz}+\delta k_z)(r_l-r_m)}+e^{-i(k_{jz}+\delta k_z)(r_l-r_m)}]\cosh^2 rS_{l,i}^+S_{m,j}^-\rho^S(t-\tau)$$

$$\approx-\frac{L}{2\pi}e^{i(\omega_i-\omega_j)t}\int_{-\infty}^\infty d\delta k_z\int_0^t d\tau e^{-i(c^2 k_{jz}\delta k_z/\omega_j)\tau}\frac{\omega_k\mu_i\mu_j}{\epsilon_0 LS\hbar}[e^{i(k_{jz}+\delta k_z)(r_l-r_m)}+e^{-i(k_{jz}+\delta k_z)(r_l-r_m)}]\cosh^2 rS_{l,i}^+S_{m,j}^-\rho^S(t-\tau)$$

$$\approx-\frac{L}{2\pi}e^{i(\omega_i-\omega_j)t}\int_0^t d\tau\frac{\omega_j\mu_i\mu_j}{\epsilon_0 LS\hbar}2\pi[e^{ik_{jz}(r_l-r_m)}\delta((r_l-r_m)-\frac{c^2 k_{jz}}{\omega_0}\tau)+e^{-ik_{jz}(r_l-r_m)}\delta((r_l-r_m)+\frac{c^2 k_{jz}}{\omega_0}\tau)]\cosh^2 rS_{l,i}^+S_{m,j}^-\rho^S(t-$$

$$\approx-\frac{L}{2\pi}e^{ik_{jz}r_{lm}}\frac{\omega_j\mu_i\mu_j}{\epsilon_0 LS\hbar}2\pi\frac{\omega_j}{c^2 k_{0z}}\cosh^2 rS_{l,i}^+S_{m,j}^-\rho^S(t)e^{i(\omega_i-\omega_j)t}$$

$$\approx-[\frac{\sqrt{\gamma_i\gamma_j}}{2}\cos(k_{0z}r_{lm})+i\frac{\sqrt{\gamma_i\gamma_j}}{2}sin(k_{0z}r_{lm})]\cosh^2 rS_{l,i}^+S_{m,j}^-\rho^S(t)e^{i(\omega_i-\omega_j)t}$$

$$(A6)$$

where emitter separation $r_{lm}=|r_l-r_m|$, collective decay rate $\gamma_i=2\mu_i^2\omega_i^2/\hbar\epsilon_0 Sc^2 k_{iz}$, and collective energy shift $\Lambda_{ij}=\sqrt{\gamma_i\gamma_j}\sin(k_{0z}r_{ij})/2$. In the third line we expand $\omega_k=c\sqrt{(\frac{\pi}{a})^2+(k_z)^2}$ around $k_z=k_{0z}$ since resonant modes provide dominant contributions. In the fifth line we extend the integration $\int_{-k_{jz}}^\infty dk_z\to\int_{-\infty}^\infty dk_z$ because the main contribution comes from the components around $\delta k_z=0$. In the next line, Weisskopf-Wigner approximation is used. Thus, we have obtained γ_{ij} and Λ_{ij} as is shown in Eq.(5).

Next we need to calculate the first term (squeezing term) in Eq.(A5):

$$e^{i(\omega_i+\omega_j-2\omega_0)t}\sum_{k_z}\int_0^t d\tau\{\boldsymbol{\mu}_{l,i}\cdot\boldsymbol{u}_{2\boldsymbol{k}_0-\boldsymbol{k}}(r_l)S_{l,i}^+\boldsymbol{\mu}_{m,j}\cdot\boldsymbol{u}_{\boldsymbol{k}}(r_m)S_{m,j}^+e^{i(\omega_{\boldsymbol{k}}-\omega_j)\tau}[-\sinh(r)\cosh(r)]\rho^S(t-\tau)$$

$$=-\frac{L}{2\pi}e^{i(\omega_i+\omega_j-2\omega_0)t}\int_0^{2k_{0z}}dk_z\int_0^t d\tau e^{i(\omega_{k_z}-\omega_j)\tau}e^{i(2k_{jz}-k_z)(r_l-o_1)}e^{ik_z(r_m-o_1)}\frac{\sqrt{\omega_{k_z}\omega_{2k_{0z}-k_z}}}{\epsilon_0 LS\hbar}\mu_i\mu_j\sinh(r)\cosh(r)S_{l,i}^+S_{m,j}^+\rho^S(t-$$

$$-\frac{L}{2\pi}e^{i(\omega_i+\omega_j-2\omega_0)t}\int_{-2k_{0z}}^0 dk_z\int_0^t d\tau e^{i(\omega_{k_z}-\omega_j)\tau}e^{i(-2k_{jz}-k_z)(r_l-o_2)}e^{ik_z(r_m-o_2)}\frac{\sqrt{\omega_{k_z}\omega_{-2k_{0z}-k_z}}}{\epsilon_0 LS\hbar}\mu_i\mu_j\sinh(r)\cosh(r)S_{l,i}^+S_{m,j}^+\rho^S(t$$

$$(A7)$$

Putting aside the overall factor $e^{i(\omega_i+\omega_j-2\omega_0)t}$, for $r_l = r_j$, Eq.(A7) reduces to

$$\sum_{k_z}\int_0^t d\tau\{\boldsymbol{\mu}_{l,i}\cdot\boldsymbol{u}_{2\boldsymbol{k}_0-\boldsymbol{k}}(r_l)S_{l,i}^+\boldsymbol{\mu}_{l,j}\cdot\boldsymbol{u}_{\boldsymbol{k}}(r_l)S_{l,j}^+ e^{i(\omega_{\boldsymbol{k}}-\omega_j)\tau}[-\sinh(r)\cosh(r)]\rho^S(t-\tau)$$

$$=-\frac{L}{2\pi}\int_0^{2k_{0z}}dk_z\int_0^t d\tau e^{i\frac{c^2 k_{jz}}{\omega_j}(k_z-k_{jz})\tau}e^{i2k_{0z}(r_l-o_1)}\frac{\sqrt{\omega_{k_z}\omega_{2k_{0z}-k_z}}\mu_i\mu_j}{\epsilon_0 LS\hbar}\sinh(r)\cosh(r)S_{l,i}^+S_{l,j}^+\rho^S(t-\tau)$$

$$-\frac{L}{2\pi}\int_{-2k_{0z}}^0 dk_z\int_0^t d\tau e^{i\frac{c^2 k_{jz}}{\omega_j}(k_z-k_{jz})\tau}e^{-i2k_{0z}(r_l-o_2)}\frac{\sqrt{\omega_{k_z}\omega_{-2k_{0z}-k_z}}\mu_i\mu_j}{\epsilon_0 LS\hbar}\sinh(r)\cosh(r)S_{l,i}^+S_{l,j}^+\rho^S(t-\tau) \qquad\text{(A8)}$$

$$=-\frac{L}{2\pi}[e^{i2k_{0z}(r_l-o_1)}+e^{-i2k_{0z}(r_l-o_2)}]\frac{\sqrt{\omega_i\omega_j}\mu_i\mu_j}{\epsilon_0 LS\hbar}\int_0^t d\tau 2\pi\delta(\frac{c^2 k_{jz}}{\omega_j}\tau)\sinh(r)\cosh(r)S_{l,i}^+S_{l,j}^+\rho^S(t-\tau)$$

$$=-e^{i2k_{jz}R}\frac{\omega_0^2\mu_i\mu_j}{\epsilon_0\hbar Sc^2 k_{0z}}\cos(2k_{0z}r_l)\sinh(r)\cosh(r)S_{l,i}^+S_{l,j}^+\rho^S(t)$$

$$=-e^{i2k_{0z}R}\frac{\sqrt{\gamma_i\gamma_j}}{2}\cos(2k_{0z}r_l)\sinh(r)\cosh(r)S_{l,i}^+S_{l,j}^+\rho^S(t)$$

where we have used the fact that the origin of coordinate system is at equal distant from two sources(i.e., $o_2 = -o_1 = R$) in the second last line. Incorporating index l into i, we have $\gamma'_{ij} = \sqrt{\gamma_i\gamma_j}\cos(2k_{0z}r_i)$. For $r_i \neq r_j$, Eq. (A7) reduces to

$$\sum_{k_z}\int_0^t d\tau\{\boldsymbol{\mu}_{l,i}\cdot\boldsymbol{u}_{2\boldsymbol{k}_0-\boldsymbol{k}}(r_l)S_{l,i}^+\boldsymbol{\mu}_{m,j}\cdot\boldsymbol{u}_{\boldsymbol{k}}(r_m)S_{m,j}^+ e^{i(\omega_{\boldsymbol{k}}-\omega_j)\tau}[-\sinh(r)\cosh(r)]\rho^S(t-\tau)$$

$$=-\frac{L}{2\pi}\int_0^{2k_{0z}}dk_z\int_0^t d\tau e^{i\frac{c^2 k_{jz}}{\omega_j}(k_z-k_{jz})\tau}e^{i2k_{0z}(r_c-o_1)}e^{-i(k_z-k_{0z})(r_l-r_m)}\frac{\sqrt{\omega_{k_z}\omega_{2k_{0z}-k_z}}\mu_i\mu_j}{\epsilon_0 LS\hbar}\sinh(r)\cosh(r)S_{l,i}^+S_{m,j}^+\rho^S(t-\tau)$$

$$-\frac{L}{2\pi}\int_{-2k_{0z}}^0 dk_z\int_0^t d\tau e^{i\frac{c^2 k_{jz}}{\omega_j}(-k_z-k_{jz})\tau}e^{-i2k_{0z}(r_c-o_2)}e^{-i(k_z+k_{0z})(r_l-r_m)}\frac{\sqrt{\omega_{k_z}\omega_{-2k_{0z}-k_z}}\mu_i\mu_j}{\epsilon_0 LS\hbar}\sinh(r)\cosh(r)S_{l,i}^+S_{m,j}^+\rho^S(t-\tau)$$

$$=-\frac{L}{2\pi}e^{i2k_{0z}(r_c-o_1)}\frac{\sqrt{\omega_i\omega_j}\mu_i\mu_j}{\epsilon_0 LS\hbar}\int_{-\infty}^\infty dk_z\int_0^t d\tau e^{i\frac{c^2 k_{jz}}{\omega_j}(k_z-k_{jz})\tau}e^{-i(k_z-k_{0z})(r_l-r_m)}\sinh(r)\cosh(r)S_{l,i}^+S_{m,j}^+\rho^S(t-\tau)$$

$$-\frac{L}{2\pi}e^{-i2k_{0z}(r_c-o_2)}\frac{\sqrt{\omega_i\omega_j}\mu_i\mu_j}{\epsilon_0 LS\hbar}\int_{-\infty}^\infty dk_z\int_0^t d\tau e^{i\frac{c^2 k_{jz}}{\omega_j}(k_z-k_{jz})\tau}e^{i(k_z-k_{0z})(r_l-r_m)}\sinh(r)\cosh(r)S_{l,i}^+S_{m,j}^+\rho^S(t-\tau)$$

$$\approx-\frac{L}{2\pi}e^{i2k_{0z}R}\frac{\omega_0^2\mu_i\mu_j}{\epsilon_0 LS\hbar}\int_0^t d\tau 2\pi[e^{i2k_{0z}r_c}\delta(r_l-r_m-\frac{c^2 k_{0z}}{\omega_0}\tau)+e^{-i2k_{0z}r_c}\delta(r_l-r_m+\frac{c^2 k_{0z}}{\omega_0}\tau)]\sinh(r)\cosh(r)S_{l,i}^+S_{m,j}^+\rho^S(t-\tau)$$

$$\approx-e^{i2k_{0z}R}\frac{\omega_0^2\mu_i\mu_j}{\epsilon_0\hbar Sc^2 k_{0z}}e^{i2k_{0z}r_c sgn(r_l-r_m)}S_{l,i}^+S_{m,j}^+\rho^S(t)\to-\frac{\sqrt{\gamma_i\gamma_j}}{2}e^{i2k_{0z}R}\cos(k_{0z}(r_l+r_m))S_{l,i}^+S_{m,j}^+\rho^S(t)$$

$$\text{(A9)}$$

where $sgn(r_l - r_m)$ is the sign function. The last arrow is because we need to sum over i,j, so the imaginary part of $e^{i2k_{0z}r_c sgn(i-j)}$ vanishes, so the neat result is that $\gamma'_{ijkl} = e^{i2k_{0z}R}\sqrt{\gamma_j\gamma_l}\cos(k_{0z}(r_i+r_k))$. As for $S_i^+\rho^S(t)S_j^+$ terms, the combination of the last two terms in Eq.(A3) will make the imaginary part of $e^{i2k_{0z}r_c sgn(r_l-r_m)}$ vanish. Thus, we have γ'_{ijkl} in Eq.(5).

APPENDIX B: DERIVATION OF COEFFICIENTS EQ.(9)

Here we will show how to derive the master equation with coefficients Eq.(B1). The mode function of the squeezed vacuum is given by

$$\boldsymbol{u}_{\boldsymbol{k}s}(\boldsymbol{r}_i) = \sqrt{\frac{\omega_{\boldsymbol{k}s}}{2\epsilon_0\hbar V}}\boldsymbol{e}_{\boldsymbol{k}s}e^{i\boldsymbol{k}\cdot(\boldsymbol{r}_i-\boldsymbol{o}_{\boldsymbol{k}s})} \qquad\text{(B1)}$$

where $\boldsymbol{o_{ks}}$ is a phenomenological parameter which includes the effects of the initial phase and the position of the squeezing source[15]. The correlation functions for the squeezed vacuum are[30]:

$$
\begin{aligned}
\left\langle a_{\boldsymbol{k},s}^{\dagger} a_{\boldsymbol{k}',s'} \right\rangle &= \sinh^2 r \delta_{\boldsymbol{k}'\boldsymbol{k}} \delta_{ss'} \\
\left\langle a_{\boldsymbol{k},s} a_{\boldsymbol{k}',s'}^{\dagger} \right\rangle &= \cosh^2 r \delta_{\boldsymbol{k}'\boldsymbol{k}} \delta_{ss'} \\
\left\langle a_{\boldsymbol{k},s}^{\dagger} a_{\boldsymbol{k}',s'}^{\dagger} \right\rangle &= -e^{-i\theta} \cosh(r)\sinh(r) \delta_{\boldsymbol{k}',2\boldsymbol{k}_0-\boldsymbol{k}} \delta_{ss'} \\
\left\langle a_{\boldsymbol{k},s} a_{\boldsymbol{k}',s'} \right\rangle &= -e^{i\theta} \cosh(r)\sinh(r) \delta_{\boldsymbol{k}',2\boldsymbol{k}_0-\boldsymbol{k}} \delta_{ss'}
\end{aligned}
\tag{B2}
$$

For simplicity, we can set the squeezing parameter $\theta = 0$, and all atoms to align along the same direction.

Since the only difference is the squeezing terms $\left\langle a_{\boldsymbol{k},s}^{\dagger} a_{\boldsymbol{k}',s'}^{\dagger} \right\rangle$, $\langle a_{\boldsymbol{k},s} a_{\boldsymbol{k}',s'} \rangle$, we will just start from Eq. (A7). Apart from the factor $e^{i(\omega_i+\omega_j-2\omega_0)t}$, When $r_i \neq r_j$, Eq. (A7) becomes:

$$
\begin{aligned}
&\sum_{k_z} \int_0^t d\tau \{ \boldsymbol{\mu}_{l,i} \cdot \boldsymbol{u}_{-2\boldsymbol{k}_0+\boldsymbol{k}}(r_l) S_{l,i}^+ \boldsymbol{\mu}_{m,j} \cdot \boldsymbol{u}_{\boldsymbol{k}}(r_m) S_{m,j}^+ e^{i(\omega_{\boldsymbol{k}}-\omega_0)\tau} [-\sinh(r)\cosh(r)] \rho^S(t-\tau) \\
&\approx -\frac{L}{2\pi} \int_0^{2k_0} dk \int_0^t d\tau e^{i\frac{c^2 k_0}{\omega_0}(k-k_0)\tau} e^{-i(2k_0-k)(r_l-o_2)+ik(r_m-o_1)} \frac{\sqrt{\omega_{k_z}\omega_{2k_{0z}-k_z}}\mu_i\mu_j}{\epsilon_0 LS\hbar} \sinh(r)\cosh(r) S_{l,i}^+ S_{m,j}^+ \rho^S(t-\tau) \\
&\quad -\frac{L}{2\pi} \int_{-2k_0}^0 dk \int_0^t d\tau e^{i\frac{c^2 k_0}{\omega_0}(-k-k_0)\tau} e^{i(2k_0+k)(r_l-o_2)+ik(r_m-o_1)} \frac{\sqrt{\omega_{k_z}\omega_{-2k_{0z}-k_z}}\mu_i\mu_j}{\epsilon_0 LS\hbar} \sinh(r)\cosh(r) S_{l,i}^+ S_{m,j}^+ \rho^S(t-\tau) \\
&= -\frac{L}{2\pi} \int_0^{2k_0} dk \int_0^t d\tau e^{i\frac{c^2 k_0}{\omega_0}(-k_0)\tau} e^{-i(2k_0)(r_l-o_2)} e^{ik(\frac{c^2 k_0}{\omega_0}\tau+r_l-o_1+r_m-o_2)} \frac{\sqrt{\omega_{k_z}\omega_{2k_{0z}-k_z}}\mu_i\mu_j}{\epsilon_0 LS\hbar} \sinh(r)\cosh(r) S_{l,i}^+ S_{m,j}^+ \rho^S(t-\tau) \\
&\quad -\frac{L}{2\pi} \int_0^{2k_0} dk \int_0^t d\tau e^{i\frac{c^2 k_0}{\omega_0}(-k_0)\tau} e^{i(2k_0)(r_l-o_2)} e^{ik(\frac{c^2 k_0}{\omega_0}\tau-r_l+o_1-r_m+o_2)} \frac{\sqrt{\omega_{k_z}\omega_{-2k_{0z}-k_z}}\mu_i\mu_j}{\epsilon_0 LS\hbar} \sinh(r)\cosh(r) S_{l,i}^+ S_{m,j}^+ \rho^S(t-\tau) \\
&\approx -\frac{L}{2\pi} \int_0^t d\tau e^{i\frac{c^2 k_0}{\omega_0}(-k_0)\tau} e^{-i(2k_0)(r_l)} e^{2ik_0 o_2} 2\pi\delta(\frac{c^2 k_0}{\omega_0}\tau+2r_c-o_1-o_2) \frac{\sqrt{\omega_i\omega_j}\mu_i\mu_j}{\epsilon_0 LS\hbar} \sinh(r)\cosh(r) S_{l,i}^+ S_{m,j}^+ \rho^S(t-\tau) \\
&\quad -\frac{L}{2\pi} \int_0^t d\tau e^{i\frac{c^2 k_0}{\omega_0}(-k_0)\tau} e^{i(2k_0)(r_l)} e^{-2ik_0 o_2} 2\pi\delta(\frac{c^2 k_0}{\omega_0}\tau-2r_c+o_1+o_2) \frac{\sqrt{\omega_i\omega_j}\mu_i\mu_j}{\epsilon_0 LS\hbar} \sinh(r)\cosh(r) S_{l,i}^+ S_{m,j}^+ \rho^S(t-\tau) \\
&\approx -\frac{L}{2\pi} e^{i(-k_0)|2r_c-o_1-o_2|} e^{sgn(r_c-o_1-o_2)i(2k_0)(r_l)} e^{-2ik_0 o_2} 2\pi \frac{\omega_0^2\mu_i\mu_j}{\epsilon_0 LS\hbar c^2 k_{0z}} \sinh(r)\cosh(r) S_{l,i}^+ S_{m,j}^+ \rho^S(t) \\
&= -\frac{L}{2\pi} e^{sgn(2r_c-o_1-o_2)ik_0(r_l-r_m)} e^{ik_0(o_1-o_2)} 2\pi \frac{\omega_0^2\mu_i\mu_j}{\epsilon_0 LS\hbar c^2 k_{0z}} \sinh(r)\cosh(r) S_{l,i}^+ S_{m,j}^+ \rho^S(t)
\end{aligned}
\tag{B1}
$$

Since we need to sum over l, m, i, j, the imaginary part of $e^{sgn(2r_c-o_1-o_2)ik_0(r_l-r_m)}$ gets canceled, which yields $\gamma'_{ijkl} = \sqrt{\gamma_i\gamma_k}\cos[k_{0z}(r_{jl})]$. The above calculation is also valid when $r_i = r_j$.

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