Generation of single entangled photon-phonon pairs via an atom-photon-phonon interaction

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Quantum blockade and entanglement play important roles in quantum information and quantum communication as quantum blockade is an effective mechanism to generate single photons (phonons) and entanglement is a crucial resource for quantum information processing. In this work, we propose a method to generate single entangled photon-phonon pairs in a hybrid optomechanical system. We show that photon blockade, phonon blockade, and photon-phonon correlation and entanglement can be observed via the atom-photon-phonon (tripartite) interaction, under the resonant atomic driving. The correlated and entangled single photons and single phonons, i.e., single entangled photon-phonon pairs, can be generated in both the weak and strong tripartite interaction regimes. Our results may have important applications in the development of highly complex quantum networks.

I. INTRODUCTION

Optomechanical systems with parametric coupling between optical and mechanical modes provide us a perfect platform for manipulating the states of photons and phonons [1]. As an important application, photon (phonon) blockade [2–4], that only allows single photon (phonon) excitation in the optical (mechanical) mode, based on optomechanical interaction has attracted significant interest in the past few years. A number of designs based on diverse mechanisms are proposed to demonstrate photon (phonon) blockade in optomechanical systems, such as photon (phonon) blockade based on strong optomechanical couplings [5–15] and photon (phonon) blockade in weak nonlinear regime induced by quantum interference [16–20].

In a recent experiment [21], the non-classical correlations between single photons and phonons from a nanomechanical resonator was reported by driving the nanomechanical photonic crystal cavity with blue-detuned optical pulse. After that, we studied the photon and phonon statistics in a quadratically coupled optomechanical system, and show that both photon blockade and phonon blockade can be observed in the same parameter regime, and more important, the single photons and single phonons are strongly anticorrelated [22]. Here, we will do a further study and propose a method to generate correlated single photons and single phonons under the constant atomic driving. Even more interestingly, we will show that the correlated single photons and single phonons are entangled with each other, i.e., they are single entangled photon-phonon pairs.

Entangled states have great significance of both fundamental physics study and applications in quantum information processing and quantum communication. The optomechanical entanglement has already been proposed theoretically [23–28] and demonstrated experimentally [29–32]. Optomechanical systems provide a perfect platform to generate both bipartite [33–36] and multipartite [37–41] entanglement. However, there are substantial differences between the entanglement we will discuss in this paper and entanglement proposed before. One striking difference is the entanglement we proposed here is for single photons and phonons, which is non-Gaussian, so that the generally adopted method, i.e., the linearization of the optomechanical interaction, is no longer applicable.

Inspired by a recent experiment [42], in which the coupling between an optomechanical resonator with two-level emitters was realized, here we consider a hybrid system which enables a tripartite interaction between a two-level atom, an optical mode, and a mechanical mode [43]. We study the generation of single entangled photon-phonon pairs, which are useful for quantum information and quantum communication. Such atom-photon-phonon interactions were proposed to provide an optically controllable interaction between a two-level atom and a macroscopic mechanical oscillator [19, 44] by driving the optical mode strongly. Nevertheless, in this paper we drive the two-level atom coherently and show that single entangled photon-phonon pairs can be generated in the hybrid optomechanical system. The single entangled photon-phonon pairs have potential application in the development of highly complex quantum networks.

The remainder of this paper is organized as follows. In Sec. II, we introduce the theoretical model of a hybrid optomechanical system, and show the simple derivation of the atom-photon-phonon interaction and the energy spectrum of the Hamiltonian. In Sec. III, the photon and

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phonon statistics, and the quantum correlation between the photons and phonons are discussed numerically. Finally, a summary is given in Sec. IV.

II. MODEL AND HAMILTONIAN

We study a hybrid system with a two-level atom (σ± being the ladder operators) of transition frequency ω0 and a mechanical resonator b of resonance frequency ωm in an optical cavity a of resonance frequency ωc, as shown in Fig. 1(a) and (b). Here, we consider a special situation in which the mechanical displacement x induces a variation of the spatial distribution of the cavity field [44], while the mechanical effect on the optical frequency ωc can be neglected. Thus, the coupling strength g(x) between the two-level atom and the optical mode depends on the position of the mechanical resonator x, which is described by the interaction Hamiltonian under the rotating-wave approximation as (ℏ = 1)

$$H_{\text{int}} = g(x)(\sigma_+ a + \sigma_- a^\dagger).$$

Typically the mechanical displacement x is very small, and g(x) can be expanded to the first order in x,

$$g(x) = g(0) + J(b^\dagger + b),$$

where $J \equiv (\partial g/\partial x)|_{x=0}$ is the tripartite atom-phonon interaction strength. In the specific condition that the two-level atom is placed at the node of the optical mode with mechanical resonator in equilibrium, i.e., $g(0) = 0$, and the only possible interaction between them is the atom-photon-phonon interaction as

$$H_{\text{int}} = J(b^\dagger + b)(\sigma_+ a + \sigma_- a^\dagger).$$

Under particular resonant conditions, the tripartite interaction allows swapping the excitation between the three quantum systems. Under the conditions $\omega_0 = \omega_c + \omega_m$, in which the mechanical displacement x induces a variation of the spatial distribution of the cavity field [44], while the mechanical effect on the optical frequency $\omega_c$ can be neglected. Thus, the coupling strength $g(x)$ between the two-level atom and the optical mode depends on the position of the mechanical resonator x, which is described by the interaction Hamiltonian under the rotating-wave approximation as ($h = 1$)

$$H_{\text{int}} = g(x)(\sigma_+ a + \sigma_- a^\dagger).$$

and min{ω0, ωc} $\gg$ $\omega_m$ $\gg$ J, the Hamiltonian of the resonant interaction reads

$$H_{\text{int}} = J(\sigma_+ ab + \sigma_- a^\dagger b^\dagger),$$

which describes the simultaneous generation of a photon and a phonon with the two-level atom jumping from the excited state to its ground state and the reverse process. This tripartite interaction provides us an effective way to generate photon-phonon pairs. Such a hybrid system can be implemented in the electromechanical systems [29, 45–47] with artificial atom at the node or in a Fabry-Pérot cavity with a membrane containing two-level atoms in the node of the cavity mode [48–51].

Next, we consider the case that the two-level atom is pumped by a coherent field (strength Ω, frequency $\omega_p$), and the total Hamiltonian for the hybrid sys-

![FIG. 1. (Color online) Schematics of the hybrid systems with atom-photon-phonon interactions: (a) a two-level atom in an optical cavity with a movable end mirror; (b) a two-level atom imbedded in a membrane inside an optical cavity. The energy spectrum of the hybrid optomechanical system [see the Hamiltonian in Eq. (5)] given (c) in the non-coupling basis and (d) in the diagonal basis.](image)
tem in the rotating frame with respect to $R(t) = \exp \left[ i \omega_p \sigma_+ t + i (\omega_p - \omega_m) a^\dagger a t + i \omega_m b^\dagger b t \right]$ reads

$$H = \Delta \sigma_+ \sigma_- + \Delta a^\dagger a + J (\sigma_+ a b + \sigma_- a^\dagger b^\dagger) + \Omega \sigma_z, \quad (5)$$

where we introduce the detuning $\Delta \equiv \omega_0 - \omega_p = \omega_c + \omega_m - \omega_p$.

The energy spectrum of the Hamiltonian in Eq. (5) for hybrid optomechanical system is shown in Figs. 1(c) and 1(d). In the non-coupling basis [Fig. 1(c)], $|e\rangle \langle g|$ denotes the excited (ground) state of the two-level atom, and $|n, m\rangle$ represents the Fock state with $n$ photons in the optical mode and $m$ phonons in the mechanical mode. In Fig. 1(d), we denote the eigenstates in the diagonal basis as $|0, 0\rangle_0 \equiv |g\rangle$, $|1, 0\rangle_0 \equiv |g\rangle|1, 0\rangle$, $|0, 1\rangle_0 \equiv |g\rangle|0, 1\rangle$, $|1, 1\rangle_0 \equiv (|g\rangle|1, 1\rangle \pm |e\rangle |0, 0\rangle)/\sqrt{2}$ with eigenvalues $0$, $\omega_c$, $\omega_m$, $\omega_0 \pm J$, respectively.

### III. CORRELATION AND ENTANGLEMENT

To quantify the statistics of the phonons and photons in the system, we consider the equal-time second-order correlation functions $g_n^{(2)}(0)$ and $g_m^{(2)}(0)$, and cross-correlation function $g_{nm}^{(2)}(0)$ in the steady state ($t \to \infty$) defined by

$$g_n^{(2)}(0) \equiv \frac{\langle a^\dagger a^\dagger a a \rangle}{\langle n \rangle^2}, \quad (6)$$

$$g_m^{(2)}(0) \equiv \frac{\langle b^\dagger b^\dagger b b \rangle}{\langle m \rangle^2}, \quad (7)$$

$$g_{nm}^{(2)}(0) \equiv \frac{\langle a^\dagger b^\dagger b a \rangle}{\langle n \rangle \langle m \rangle}, \quad (8)$$

where $\langle n \rangle \equiv \langle a^\dagger a \rangle$ and $\langle m \rangle \equiv \langle b^\dagger b \rangle$ are the mean photon and phonon numbers. The dynamic behavior of the total open system is described by the master equation for the density matrix $\rho$ of the system [52]

$$\frac{\partial \rho}{\partial t} = -i [H, \rho] + \kappa L[\sigma_-] \rho + \gamma_c L[a] \rho + \gamma_m (m_{th} + 1) L[b] \rho + \gamma_m m_{th} L[b^\dagger] \rho, \quad (9)$$

where $L[\sigma_-] \rho = opo^\dagger - (o^\dagger \rho o + \rho o^\dagger o) / 2$ denotes a Lindblad term for an operator $o$; $\kappa$ is the damping rate of the two-level atom and $\gamma_c$ is the damping rate of the optical (mechanical) mode; $m_{th}$ is the mean thermal phonon number. We assume that the frequencies of the two-level atom and the optical mode are so high that the thermal effect can be neglected.

The equal-time second-order correlation functions $[g_n^{(2)}(0)$ and $g_m^{(2)}(0)]$ and cross-correlation function $g_{nm}^{(2)}(0)$ are plotted as functions of the detuning $\Delta/\kappa$ in Fig. 2 under both weak-coupling condition $||a\rangle J = 0.1\kappa$ and strong-coupling condition $||c\rangle J = 100\kappa$. It is clear that photon blockade and phonon blockade, i.e., $g_n^{(2)}(0) = g_m^{(2)}(0) < 1$, appear simultaneously around $|\Delta| = J$. 

FIG. 2. (Color online) In panels (a) and (d), the equal-time second-order correlation functions $[g_n^{(2)}(0)$ and $g_m^{(2)}(0)]$ and cross-correlation function $g_{nm}^{(2)}(0)$ are plotted as functions of the detuning $\Delta/\kappa$. In panels (b) and (e), the mean photon (phonon) number $|\langle n \rangle = |\langle m \rangle|$ is plotted as a function of the detuning $\Delta/\kappa$. In panels (c) and (f), the logarithmic negativity $E_N$ is plotted as a function of the detuning $\Delta/\kappa$. We set $J = 0.1\kappa$ in (a) and (b) and set $J = 100\kappa$ in (c) and (d). Other used parameters are $\gamma_c = \gamma_m = 10\kappa$, $\Omega = \kappa$, and $m_{th} = 0$. 

In panels (a) and (b) of Fig. 2, we denote the eigenstates in the diagonal basis $|0, 0\rangle_0 \equiv |g\rangle$, $|1, 0\rangle_0 \equiv |g\rangle|1, 0\rangle$, $|0, 1\rangle_0 \equiv |g\rangle|0, 1\rangle$, $|1, 1\rangle_0 \equiv (|g\rangle|1, 1\rangle \pm |e\rangle |0, 0\rangle)/\sqrt{2}$ with eigenvalues $0$, $\omega_c$, $\omega_m$, $\omega_0 \pm J$, respectively.
the same parameters. Simultaneously, the single photons and single phonons generated by photon blockade and phonon blockade are strongly correlated with each other, i.e., $g_{nm}^{(2)}(0) \gg 1$. The optimal detuning $\Delta$ for correlated photon blockade and phonon blockade depends on the coupling strength $J$: $\Delta = 0$ for weak coupling and $|\Delta| \approx J$ for strong coupling. Moreover, the mean photon (phonon) number $\langle n \rangle = \langle m \rangle$ in the weak-coupling case is much smaller than the one in the strong-coupling case.

Physically, the single photon and phonon pairs are generated one by one with the two-level atom jumping from the excited state to its ground state. In the weak-coupling regime ($J \ll \kappa$), the system is driven resonantly with detuning $\Delta = 0$ because the states $|1, 1\rangle_+ \text{ and } |1, 1\rangle_-$ are not resolved. In the strong-coupling regime ($J \gg \kappa$), the system should be investigated by the dressed states as shown in Fig. 1(d), and the single photon and phonon pairs are generated with detuning $\Delta = \pm J$ for resonant pumping.

In order to understand the behavior of the cross-correlation function $g_{nm}^{(2)}(0)$, we can give the expression of $g_{nm}^{(2)}(0)$ approximately. Under the weak-exciting condition, i.e., $\max\{\langle n \rangle, \langle m \rangle\} \ll 1$, we have mean photon (phonon) number

\begin{equation}
\langle n \rangle \approx \rho_{5,5} + \rho_{4,4},
\end{equation}

\begin{equation}
\langle m \rangle \approx \rho_{5,5} + \rho_{3,3},
\end{equation}

and the cross-correlation function $g_{nm}^{(2)}(0)$

\begin{equation}
g_{nm}^{(2)}(0) \approx \frac{\rho_{5,5}}{(\rho_{5,5} + \rho_{3,3})(\rho_{5,5} + \rho_{4,4})},
\end{equation}

where $\rho_{3,3} = \langle g\langle 0, 1|\rho|g\rangle\langle 0, 1 \rangle \rangle$, $\rho_{4,4} = \langle g\langle 1, 0|\rho|g\rangle\langle 1, 0 \rangle \rangle$, and $\rho_{5,5} = \langle g\langle 1, 1|\rho|g\rangle\langle 1, 1 \rangle \rangle$, and they satisfy the relations

\begin{equation}
\rho_{3,3} \approx \frac{\gamma_c}{\gamma_m} \rho_{5,5}.
\end{equation}

\begin{equation}
\rho_{4,4} \approx \frac{\gamma_m}{\gamma_c} \rho_{5,5}.
\end{equation}

If we set $\gamma_c = \gamma_m$, then we have $\rho_{3,3} \approx \rho_{4,4} \approx \rho_{5,5}$, $\langle n \rangle = \langle m \rangle \approx 2\rho_{5,5}$, and

\begin{equation}
g_{nm}^{(2)}(0) \approx \frac{1}{2\langle n \rangle}.
\end{equation}

Under the resonant conditions at the detuning $|\Delta| = J$, we have maximum $\langle n \rangle = \langle m \rangle$, and thus minimum cross-correlation function $g_{nm}^{(2)}(0)$, corresponding to the dips around the detuning $|\Delta| = J$.

It’s not hard to guess that the strongly correlated single photons and single phonons generated by photon blockade and phonon blockade are entangled with each other. The entanglement between the optical and mechanical modes can be characterized by the logarithmic negativity [53]

\begin{equation}
E_N = \log_2 \left\| \rho_{AB}^{T_A} \right\|_1,
\end{equation}

where the symbol $\|\cdot\|_1$ denotes the trace norm, and $\rho_{AB}^{T_A}$
is the partial transpose of the reduced density matrix $\rho_{AB}$ of the optical and mechanical modes. It is worth mentioning that the entangled state for the single photons and single phonons obtained here is non-Gaussian. Thus the logarithmic negativity for Gaussian states [59] widely used in the previous works [23–28] cannot be used to accurately describe the entangled state here.

The logarithmic negativity $E_N$ is shown in Figs. 2(c) and 2(f). Obviously, the strongly correlated single photons and single phonons generated by photon blockade and phonon blockade are entangled with each other in both the weak ($J < \kappa$) and strong ($J > \kappa$) coupling regimes. In the weak-coupling regime as shown in Fig. 2(c), there is a dip around the detuning $\Delta = 0$, which is induced by the quantum interferences between two routes: (a) the direct transition channel $|g\rangle |0,0\rangle \xrightarrow{\Omega} |e\rangle |0,0\rangle \xrightarrow{\gamma} |g\rangle |1,1\rangle$; (b) the indirect transition channel $|g\rangle |0,0\rangle \xrightarrow{\Omega} |e\rangle |0,0\rangle \xrightarrow{\Omega} |e\rangle |0,0\rangle \xrightarrow{\gamma} |g\rangle |1,1\rangle$ (or higher-order variants). Thus the width of the dip depends on the driving strength $\Omega$, as shown in Fig. 3(a). Similar mechanism can induce transparency in lambda-type three-level atoms [54, 55] and optomechanical systems [56–58]. Differently, in Fig. 2(f), there are two peaks around the detunings $\Delta = \pm J$ in the strong-coupling regime. This phenomenon can be understood by analyzing the energy spectrum shown in Fig. 1(d): the transition process $|0,0\rangle \rightarrow |1,1\rangle$ is resonantly enhanced with detunings $\Delta = \pm J$. As a consequence, we can shift the optimal value of the detuning for entanglement by tuning the coupling strength $J$ as shown in Fig. 3(b).

Figure 4 shows the second-order correlation functions $g_n^{(2)}(0)$ and $g_m^{(2)}(0)$ and cross-correlation function $g_m(0)$ with the coupling strength $J$ from weak to strong. The mean photon (phonon) number $|\langle n \rangle| = |\langle m \rangle|$ and logarithmic negativity $E_N$ increase with the enhancing of the coupling strength $J$. As shown in Fig. 4(b), the cross-correlation function $g_m^{(2)}(0)$ decreases with the increasing of the mean photon (phonon) number $|\langle n \rangle| = |\langle m \rangle|$, and the numerical results (red dashed curve) agrees well with the analytical results given in Eq. (15) (blue short-dashed curve). The second-order correlation functions $g_n^{(2)}(0)$ and $g_m^{(2)}(0)$ increases first with the mean photon (phonon) number, and then decreases with the coupling strength $J$, as the excitations of states $|2,1\rangle_{\pm} = (|g\rangle |2,1\rangle \pm |e\rangle |1,0\rangle) / \sqrt{2}$ are suppressed for the enhancement of the effective damping rates with the coupling strength $J$. The suitable coupling strength $J$ for observing correlated single photons and single phonons, i.e., $g_n^{(2)}(0) = g_m^{(2)}(0) \ll 1$ and $g_m^{(2)}(0) \gg 1$, is $J < \kappa$ or $J \gg \kappa$.

Generally, the damping rate of the mechanical mode is much smaller than the damping rate of the optical mode. However, the effective damping of the mechanical mode can be controlled and significantly enhanced by coupling the mechanical mode to an auxiliary opt-

![FIG. 5. (Color online) (a) and (c) the equal-time second-order correlation functions $g_n^{(2)}(0)$ and $g_m^{(2)}(0)$ and cross-correlation function $g_m^{(2)}(0)$ are plotted as functions of the mechanical damping rate $\log_{10}[\gamma_m / \kappa]$. (b) and (d), the mean photon (phonon) number $|\langle n \rangle|$ and $|\langle m \rangle|$ and the logarithmic negativity $E_N$ are plotted as functions of the mechanical damping rate $\log_{10}[\gamma_m / \kappa]$. We set $J = 0.1\kappa$ in (a) and (b) and set $J = 100\kappa$ in (c) and (d). Other used parameters are $|\Delta| = J$, $\gamma_c = 10\kappa$, $\Omega = \kappa$, and $\kappa_{th} = 0$.](image)

![FIG. 6. (Color online) The elements of the density matrix $\rho$ from Eq. (9) in the steady state are plotted as functions of the mechanical damping rate $\log_{10}[\gamma_m / \kappa]$, where $\rho_{1,1} = |g\rangle \langle 0,0| \rho |g\rangle \langle 0,0|$, $\rho_{2,2} = |e\rangle \langle 0,0| \rho |e\rangle \langle 0,0|$, $\rho_{3,3} = |g\rangle \langle 0,1| \rho |g\rangle \langle 0,1|$, $\rho_{4,4} = |g\rangle \langle 1,0| \rho |g\rangle \langle 1,0|$, $\rho_{5,5} = |g\rangle \langle 1,1| \rho |g\rangle \langle 1,1|$, $\rho_{6,6} = |g\rangle \langle 0,0| \rho |g\rangle \langle 1,1|$, $\rho_{7,7} = |e\rangle \langle 0,0| \rho |e\rangle \langle 1,1|$, $\rho_{8,8} = |e\rangle \langle 0,0| \rho |e\rangle \langle 1,1|$. We set $J = 0.1\kappa$ in (a) and set $J = 100\kappa$ in (b). Other used parameters are $|\Delta| = J$, $\gamma_c = 10\kappa$, $\Omega = \kappa$, and $\kappa_{th} = 0$.](image)
FIG. 7.  (Color online) (a) and (c) the equal-time second-order correlation functions $[g_n^{(2)}(0)$ and $g_m^{(2)}(0)]$ and cross-correlation function $g_m^{(2)}(0)$ are plotted as functions of the mean thermal phonon number $\log_{10}[m_{th}]$.  (b) and (d), the mean photon (phonon) number $\langle n \rangle$ and $\langle m \rangle$ and the logarithmic negativity $E_N$ are plotted as functions of the mean thermal phonon number $\log_{10}[m_{th}]$.  We set $J = 0.1k$ in (a) and (b) and set $J = 100k$ in (c) and (d).  Other used parameters are $|\Delta| = J$, $\gamma_c = \gamma_m = 10k$, and $\Omega = k$.

In conclusion, we have studied the photon and phonon statistics, and the quantum correlation between photons and phonons in a hybrid optomechanical system including an atom-photon-phonon (tripartite) interaction.  We have shown that both the photon and phonon blockade can be observed in the same parameter area, and the generated single photons and single phonons are correlated and entangled with each other.  Moreover, the single entangled photon-phonon pairs can be observed in both the weak and strong tripartite interaction regime.  The phonons with low-loss can be used for quantum memories, and photons are suitable for the transmission of quantum information.  The generated single entangled photon-phonon pairs will have applications in quantum communication and the hybrid optomechanical system can serve as quantum transducers in building hybrid quantum networks.  In addition, the basic mechanism of this work can be generalized to a nondegenerate two-photon Jaynes-Cummings model [68, 69], to generate entangled photon pairs with different frequency, such as entangled microwave-optical photon pairs [70].
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