Detecting quantum phase transitions of photons through a defect cavity

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New Journal of Physics 13 (2011) 083036 (11pp)
Received 9 April 2011
Published 31 August 2011
Online at http://www.njp.org/
doi:10.1088/1367-2630/13/8/083036

Abstract. The quantum phase transition of photons between the Mott-insulating and superfluid states can take place in a periodic array of nanocavities, where the photons are strongly coupled to a two-level system inside each cavity. Here we consider a defect cavity inside the array structure and examine the dynamics of the corresponding photons using the non-equilibrium formalism. The global signatures of the phase transition, such as the compressibility and order parameter, are revealed in the defect-cavity photon number. Furthermore, an open electron transport device (a transport qubit) can be embedded in the defect cavity for the detection of the quantum phase transition.

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1. Introduction

The quantum phase transition has been widely studied to understand the quantum nature of many-body systems at zero temperature and to manipulate the ground states through tuning the parameters of the system Hamiltonian [1]. For interacting bosons, the celebrated Bose–Hubbard model was proposed to exhibit a transition between the Mott-insulating (MI) and superfluid (SF) phases [2]. Such a phase transition has been achieved by using ultracold atoms held inside an optical lattice, subsequent to the observation of their Bose–Einstein condensation (BEC), which corresponds to the SF phase [3, 4].

The Bose–Hubbard model is not applicable to photons since there is no photon–photon interaction. However, the MI–SF transition of photon states could take place in the strong-coupling regime of the cavity-quantum-electrodynamics systems. In this regime, the Jaynes–Cummings–Hubbard (JCH) model is used to describe the Hamiltonian of a cavity-array structure, where photons effectively interact with each other through the coupling to individual two-level systems [5–12]. Such a system can be realized by using photonic-crystal nanocavities with quantum-dot excitons [13–15], microwave cavities with Cooper-pair charge qubits [16, 17] or slot-waveguide cavities with the nitrogen-vacancy centers of diamond [18–20]. Compared with ultracold atoms, a benefit of the cavity-array systems is that one could manipulate each site of a large lattice [21]. Another example is a controllable scheme for polaritons inside a cavity-array structure, which is proposed to achieve the MI–SF transition in the Bose–Hubbard model and a phase of highly entangled and delocalized states with attractive on-site potentials [22].

To realize the photonic phase transition, one important issue is measuring the global quantities of the macroscopic states of photons, such as the compressibility and the order parameter. In this paper, we study how a weakly coupled single defect inside a JCH lattice reflects the global signatures of the lattice, and propose a scheme to detect the transition through this defect by using electron transport techniques. The defect here means that one of the cavities contains a particular dot (qubit) different from other cavities in the lattice and its inter-cavity coupling is much weaker than that of others. For a large enough JCH lattice, the defect reasonably has minor affection to the lattice. The defect cavity is found to provide enough information on the phase transition, which can be detected by coupling the cavity to a non-equilibrium electron transport device (figure 1), i.e. a transport qubit (TQ) [23–26].

2. The Jaynes–Cummings–Hubbard lattice

The Hamiltonian of the JCH lattice is (setting $\hbar = 1$ throughout)

$$H_p = \sum_i \left[ \frac{\epsilon}{2} \sigma^z_i + \omega a_i^\dagger a_i + \beta (\sigma^{+}_i a_i + \sigma^{-}_i a_i^\dagger) \right] - \kappa \sum_{\langle i, j \rangle} a_i^\dagger a_j,$$

where $i$ and $j$ represent the lattice sites, $\epsilon$ is the energy-level splitting, $\omega$ is the cavity-photon energy, $\beta$ is the coupling between the dot and photons and $\kappa$ is the inter-cavity tunneling between two nearest-neighbor sites. The two-level system of the $i$th site is described by the Pauli matrices $\sigma^z_i$ and $\sigma^{\pm}_i$, and $a_i$ is the annihilating operator of the $i$th cavity photons. The zero-temperature ground state is determined by the generalized Hamiltonian [5, 7–12]

$$\mathcal{H}_p = H_p - \mu \sum_i (a_i^\dagger a_i + \sigma^{+}_i \sigma^{-}_i),$$

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Figure 1. A diagram of a hexagonal cavity–dot (red circles filled with black or light-blue colors) array structure based on a photonic crystal. Each black site is a cavity containing a two-level quantum dot with level spacing $\epsilon$ and cavity–dot coupling $\beta$ (lower right panel), and these sites form a lattice or a polariton reservoir with inter-cavity photon tunneling $\kappa$. The light-blue site (defect) consists of a cavity and a two-level QD, $|L\rangle$ and $|R\rangle$, with energy spacing $\epsilon_s$ (lower left panel). An electron tunnels from the source to $|L\rangle$, and oscillates between $|L\rangle$ and $|R\rangle$ through stimulation by a coherently driving laser ($\beta_{co}$) and coupling to the defect-cavity photons ($\beta_s$). It then tunnels from $|R\rangle$ to the drain. The electron tunneling rates of the source and drain are $\Gamma_L$ and $\Gamma_R$, respectively. The photon tunneling rate $\kappa_s$ between the defect site and the polariton reservoir is assumed to be much weaker than $\kappa$ such that the dynamics of the defect photons is Markovian.

where $\mu$ is the chemical potential. It can be solved by introducing an order parameter $\psi \equiv \langle a_i \rangle$ and making a mean-field decoupling between cavities, i.e. $a_i^\dagger a_j \approx \langle a_i^\dagger \rangle a_j + a_j^\dagger \langle a_i \rangle - \langle a_i^\dagger \rangle \langle a_j \rangle$, by ignoring the second order of the amplitude fluctuation [5, 10]. In an experiment, a pumping laser is necessary to overcome the loss of cavity photons and decay of quantum dots [6, 27–30]. If the whole system reaches thermal equilibrium much faster than the processes of pumping and decay, one can effectively tune the chemical potential by controlling the pumping laser. Recent experiments have made significant progress to produce a photonic chemical potential and create the photonic BEC [29, 30]. In the strong-coupling regime, the photonic phase transition can be manipulated by tuning the chemical potential, as shown in figure 2. In the MI phase, the number of polaritons is certain, and the order parameter is zero. The average polariton number per cavity, $n_{po} \equiv \langle a_i^\dagger a_i + \sigma_i^+ \sigma_i^- \rangle$, exhibits integer plateaus, i.e. the polaritons are incompressible in this phase. Otherwise the system is in the SF phase. In addition, the average photon number $n_{ph} = \langle a_i^\dagger a_i \rangle$ also characterizes the phase transition although it is not a conserved quantity. The global compressibility of photons can be defined as

$$C_{ph} \equiv \frac{\partial n_{ph}}{\partial \mu},$$

which is a part of the polariton compressibility.
Figure 2. (a) The order parameter $\psi$, (b) the average polariton number and (c) the average photon number of the JCH lattice with detuning $\Delta \equiv \omega - \epsilon = 0$. The MI phases exhibit a zero-order parameter with integer polariton numbers and half-integer photon numbers, while the SF phases correspond to a non-zero-order parameter.

3. The defect cavity and transport qubit

In order to detect the Mott transition, we consider a defect cavity inside the lattice, with a TQ embedded in the defect. The defect cavity and the TQ are considered as our ‘system’, while the surrounding JCH lattice represents the ‘polariton reservoir’ (figure 1). The system Hamiltonian is

$$H_s \equiv \frac{\epsilon_s}{2} \left( c_R^\dagger c_L - c_L^\dagger c_R \right) + \beta_{c0} (c_R^\dagger c_R + c_L^\dagger c_L) + \omega a_s^\dagger a_s + \beta_a (c_R^\dagger c_R a_s + c_L^\dagger c_L a_s^\dagger).$$

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The ground-state signatures of the polariton reservoir are substantially reflected in $V$ and $\Upsilon$; the uncertainty results in the appearance of like Hamiltonian [\[\text{...}\]] or the tight-binding model [\[\text{...}\]]. Together with the rotating-wave approximation, we have

$$\mathcal{L}_p[\rho_s(t)] = -i[V^{MF}, \rho_s(0)] - \Upsilon_+[[a_s, a_s^\dagger \rho_s(t)] - [a_s^\dagger, \rho_s(t)a_s]]$$

$$- \Upsilon_-[[a_s^\dagger, a_s \rho_s(t)] - [a_s, \rho_s(t)a_s^\dagger]],$$

where the initial state $\rho_s(0)$ is decoupled from the reservoir and satisfies $[\rho_s(0), H_s] = 0$ and $V^{MF} \equiv \text{Tr}_p(V) = -\kappa_s \psi(a_s + a_s^\dagger)$. The photon tunneling rates are $\Upsilon_+ = \Upsilon_0[z\kappa_s^2/(z-1)|\psi|^2]$ and $\Upsilon_- = \Upsilon_+ + \Upsilon_0$, where $\Upsilon_0 = 2\pi z\kappa_s^2/W$ is independent of the state of the polariton reservoir. The ground-state signatures of the polariton reservoir are substantially reflected in $V^{MF}$ and $\Upsilon_\pm$. However, unlike the usual assumption of particle-number certainty of the reservoir, the polariton number is uncertain in the SF phase due to the spontaneous symmetry breaking, and the uncertainty results in the appearance of $\psi$ in the master equation.
4. Dynamics of defect photons

In the absence of the TQ, the time evolution of the average defect photon number and amplitude is given by

$$\frac{d}{dt}\langle a_t^\dagger a_s \rangle_t = -2\gamma_0 \langle a_t^\dagger a_s \rangle_t + iz\kappa_s \psi \langle a_t^\dagger - a_s \rangle_0 + 2\gamma_+,$$

$$\frac{d}{dt}\langle a_s \rangle_t = (-i\omega - \gamma_0) \langle a_s \rangle_t + iz\kappa_s \psi.$$

There is an analytic solution of (9). Assuming that the initial state is the vacuum state, it reaches

$$\langle a_t^\dagger a_s \rangle_{t \to \infty} = n_{ph} + (z - 1) |\psi|^2 = \frac{\gamma_+}{\gamma_0},$$

$$\langle a_s \rangle_{t \to \infty} = \frac{iz\kappa_s \psi}{i\omega + \gamma_0}$$

for the stationary state (dashed lines in figure 3(a)). The stationary defect-photons state characterizes the quantum phase transition of the reservoir. For MI phases, the photon number exhibits the same plateaus as the reservoir, and the amplitude (proportional to $|\psi|$) is zero. In the SF phase, the photon amplitude is reduced, but the leakage of photons from the reservoir to the defect is particularly enhanced due to the spatial coherence of the reservoir photon field.

5. Electron tunneling current

The electron flow operator (or the current) through the TQ is defined as $J \equiv -i[c_R^\dagger R c_L c_L^\dagger c_R - c_L^\dagger c_L c_R c_L]$, which consists of a coherent part $J_{co} = -i\beta_{co}(c_R^\dagger R c_L - c_L^\dagger c_R)$ and an incoherent part $J_s = -i\beta_s(c_R^\dagger R c_L a_s - c_L^\dagger c_R a_s)$. In the off-resonant regime ($\Delta_s \equiv \omega - \epsilon_s \gg \beta_s$), the coupling between the TQ and defect cavity becomes an ac Stark shift $-\beta_s^2/(\Delta_s)(c_L^\dagger c_L - c_R^\dagger c_R)a_t^\dagger a_s$ [46, 47], and $J_s$ can be ignored. With the approximation $a_t^\dagger a_s \approx \langle a_t^\dagger a_s \rangle_{t \to \infty}$, the TQ is decoupled from photons, and its equations of motion are determined by

$$\frac{d}{dt}\langle c_L^\dagger c_R \rangle_t = (\epsilon_s' - \Gamma_R) \langle c_L^\dagger c_R \rangle_t - i\beta_{co}(\langle c_L^\dagger c_L \rangle_t - \langle c_R^\dagger c_R \rangle_t),$$

$$\frac{d}{dt}\langle c_R^\dagger c_L \rangle_t = (-i\epsilon_s' + \Gamma_R) \langle c_R^\dagger c_L \rangle_t + i\beta_{co}(\langle c_R^\dagger c_R \rangle_t - \langle c_L^\dagger c_L \rangle_t),$$

$$\frac{d}{dt}\langle c_L^\dagger c_L \rangle_t = 2\Gamma_L - 2\Gamma_L \langle c_L^\dagger c_L \rangle_t - 2\Gamma_L \langle c_R^\dagger c_R \rangle_t - i\beta_{co}(\langle c_L^\dagger c_L \rangle_t - \langle c_R^\dagger c_R \rangle_t),$$

$$\frac{d}{dt}\langle c_R^\dagger c_R \rangle_t = -2\Gamma_R \langle c_R^\dagger c_R \rangle_t + i\beta_{co}(\langle c_L^\dagger c_L \rangle_t - \langle c_R^\dagger c_R \rangle_t),$$

where $\epsilon_s' \equiv \epsilon_s - (\beta_s^2/\Delta_s)(2\langle a_t^\dagger a_s \rangle_{t \to \infty} + 1)$ is just the energy difference between the Stark/Lamb-shifted $|L\rangle$ and $|R\rangle$ states.

The steady-state current is given by

$$\langle J_{co} \rangle_{t \to \infty} = \frac{2\beta_{co}^2 \Gamma_R}{\beta_{co}^2(2 + \Gamma_R/\Gamma_L) + \Gamma_R^2 + \epsilon_s'^2},$$

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Figure 3. (a) The average photon number of the defect cavity with (solid lines) and without (dashed lines) the back action from the resonant TQ. (b) The off-resonant ($\Delta_s = 0.999$) electron current with $\beta_{co} = 0.01$ and $\Gamma_L = \Gamma_R = 0.1$ and (c) the resonant ($\Delta_s = 0$) current through the TQ with $\beta_{co} = 0$ and $\Gamma_L = \Gamma_R = 0.0001$. Other parameters are $\beta_s = 0.1$, $\omega = 1$ and $\kappa_s = 10^{-4}$, and all parameters are in units of $\beta$. The quantum phase transition of the reservoir is totally reflected in the defect cavity and the TQ, where the plateaus correspond to the MI phases with integer polariton number $n_{po}$ of the reservoir and otherwise the reservoir is in the SF phase.

where the initial state of the TQ is set to $|G\rangle$. The defect-photon signatures are detected through the shifted energy spacing $\epsilon_s'$, which is proportional to the photon number. In the off-resonant limit, the defect photon number also obeys the same equation of motion in (9), and thus an ideally non-destructive passive measurement is achieved without back action to the defect.
cavity. From figure 3(b), the steady-state current exhibits plateaus for MI phases and significant dips for SF phases. The current is independent of $\kappa$ but only depends on the defect photon number. This leads to a one-to-one and inverse-square correspondence between the plateaus of the photon number and current.

In the resonant limit ($\Delta_s = 0$), the coherent tunneling is turned off ($\beta_{co} = 0$) for simplicity, and the electron current only comprises the incoherent part $J_s$ accompanying stimulated emission of photons. The equations of motion of the TQ are

$$\frac{d}{dt} \langle c_L^\dagger c_L \rangle_t = -(J_s)_t + 2\Gamma_L (1 - \langle c_L^\dagger c_L \rangle_t - \langle c_R c_R \rangle_t),$$

$$\frac{d}{dt} \langle c_R^\dagger c_R \rangle_t = +(J_s)_t - 2\Gamma_R \langle c_R^\dagger c_R \rangle_t,$$

$$\frac{d}{dt} \langle J_s \rangle_t = 2\beta_s^2 [(\langle c_L^\dagger c_L - c_R^\dagger c_R \rangle a_s^\dagger a_s)_t + \langle c_L^\dagger c_L \rangle_t] - (\Gamma_R + \Upsilon_0) \langle J_s \rangle_t,$$

and the defect cavity is strongly coupled to the TQ with

$$\frac{d}{dt} \langle a_s^\dagger a_s \rangle_t = (J_s)_t + 2\Upsilon_0 - 2\Upsilon_0 \langle a_s^\dagger a_s \rangle_t.$$

To get an analytic solution, we make the approximation

$$\langle (c_L^\dagger c_L - c_R^\dagger c_R a_s^\dagger a_s)_t \rangle_t \approx \langle c_L^\dagger c_L - c_R^\dagger c_R \rangle_t \langle a_s^\dagger a_s \rangle_t,$$

(15)

to (13), and furthermore $\langle a_s^\dagger a_s \rangle_t$ can be replaced by $\langle a_s^\dagger a_s \rangle_{t\to\infty}$ as an input parameter. Therefore, the steady-state solution is obtained in a self-consistent manner and is given by

$$\langle a_s^\dagger a_s \rangle_{t\to\infty} = \frac{(J_s)_{t\to\infty} + 2\Upsilon_0}{2\Upsilon_0},$$

$$\langle J_s \rangle_{t\to\infty} = J_1 + \sqrt{J_1^2 + \frac{4\Gamma_R \Upsilon_-}{2 + \Gamma_R/\Gamma_L}},$$

(16)

where

$$J_1 = \left(1 - \frac{\Upsilon_0 \Gamma_R}{\beta_s^2} \right) \frac{\Upsilon_0 + \Gamma_R}{2 + \Gamma_R/\Gamma_L} - \Upsilon_-.$$

(17)

From (16), the TQ has strong back action to the defect cavity. When the current is large, the TQ emits a large number of photons and the phase-transition signature is obscure. However, the back action can be suppressed by reducing the electron tunneling rates to the order of $\kappa_s/\kappa$. The resonant current is plotted in figure 3(c), and the solid lines of figure 3(a) represent the back action to the defect photons. With weak tunneling rates of the TQ, the current essentially exhibits the phase transition. In contrast with the off-resonant case, the current value reduces and the back action grows with $\kappa$, resulting from that the reservoir with a larger $\kappa$ compensates weakly for the back action ($\Upsilon_0 \propto 1/\kappa$). The $\kappa$ dependence does not depreciate the resonant TQ since $\kappa$ is not actually tunable but determined during the cavity fabrication. In addition, the absence of the coherent tunneling even makes the scheme easier to achieve in experiments.
6. Discussion

If one tries to measure the photon states of a cavity site of the reservoir, the results only represent the local properties, especially for the presence of disorders. Despite that the measurement itself could alter the photon states, it is difficult to get the ensemble-average measurements of all sites. The interference pattern of the leaking light among different cavities can be used to identify the SF phase based on the spatial coherence. However, this method does not provide enough information to distinguish or identify the MI phases with zero-order parameters, since losing coherence cannot guarantee the certainty of particle number.

In our scheme, the defect cavity provides a measure of two global properties of the reservoir simultaneously, i.e. the photon compressibility and the order parameter, through the local compressibility of the defect photons

$$C_d \equiv C_{\text{ph}} + (z - 1) \frac{\partial}{\partial \mu} |\psi|^2. \quad (18)$$

This is a direct result of the defect photon number (10). Due to the weak coupling $\kappa_s$, the photon tunneling (in and out) processes of the defect cavity are much slower than the thermalization of the reservoir, and the tunneling rates ($\Upsilon \propto n_{\text{ph}} + (z - 1) |\psi|^2$) determine the final state of the defect photons. For the SF phase, the coherence and delocalization of the reservoir photons enhance $\Upsilon$, and the reservoir order parameter ultimately contributes to extra accumulation of the defect photons.

The general relation between the derivative of TQ currents and the reservoir properties can be derived by differentiating (12) and (16) with respect to $\mu$,

$$\frac{\partial}{\partial \mu} \langle J \rangle_{t \to \infty} = A_1 C_d = A_1 \left[ C_{\text{ph}} + (z - 1) \frac{\partial}{\partial \mu} |\psi|^2 \right], \quad (19)$$

where $A_1$ is a function of $n_{\text{ph}}$, $\psi$ and defect parameters for the off-resonant case and it has an additional dependence on $\kappa$ for the resonant case. The first-order current derivative is zero for MI phases and for the local minima of dips in SF phases, owing to the contribution of the order parameter in (19). The MI phases can be easily distinguished from the local minima by higher-order derivatives.

It is inevitable to produce some disorder during the fabrication of the photonic-crystal-cavity arrays, such as fluctuations of on-site energies and inter-cavity coupling. Artificial disorders could also be introduced by altering the number of quantum dots inside each cavity. Such disorders could lead to glass phases with finite compressibility and no phase coherence of the polariton reservoir [2, 48, 49]. When a glass phase forms between two MI phases, from (19), the TQ current varies smoothly without a dip between two plateaus due to the loss of coherence. The transition from SF to glass phases could also be identified by the difference in the current derivatives.

7. Summary

In conclusion, we have considered a defect cavity coupled weakly to the Jaynes-Cummings–Hubbard lattice, and proposed a scheme using an electron TQ coupled to the defect to measure the quantum phase transition of the lattice. The dynamics of the photon tunneling between the defect and lattice is studied, and the analytic solutions of the stationary states are obtained. The defect cavity is found to retrieve global information on the lattice.
photon in equilibrium and plays the role of an interface between the TQ and the lattice. For either the resonant or off-resonant regime, the steady-state current contains information on the compressibility and order parameter, which can be used to identify different Mott insulators, the transition to superfluidity or even the glass phase resulting from disorders. Our scheme can also be extended to other coupled-cavity-array systems such as strongly correlated exciton polaritons.

Acknowledgment

This work is partially supported by the National Science Council, Taiwan, under the grant number NSC 98-2112-M-006-002 MY3.

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