A photonic Carnot engine powered by a spin-star network

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received 22 November 2016; accepted in final form 6 April 2017
published online 26 April 2017

PACS 03.67.-a – Quantum information
PACS 05.70.-a – Thermodynamics
PACS 03.65.Yz – Decoherence; open systems; quantum statistical methods

Abstract – We propose a spin-star network, where a central spin-(1/2), acting as a quantum fuel, is coupled to N outer spin-(1/2) particles. If the network is in thermal equilibrium with a heat bath, the central spin can have an effective temperature, higher than that of the bath, scaling nonlinearly with N. Such temperature can be tuned with the anisotropy parameter of the coupling. Using a beam of such central spins to pump a micromaser cavity, we determine the dynamics of the cavity field using a coarse-grained master equation. We find that the central-spin beam effectively acts as a hot reservoir to the cavity field and brings it to a thermal steady state whose temperature benefits from the same nonlinear enhancement with N and results in a highly efficient photonic Carnot engine. The validity of our conclusions is tested against the presence of atomic and cavity damping using a microscopic master equation method for typical microwave cavity-QED parameters. The role played by quantum coherence and correlations on the scaling effect is pointed out. An alternative scheme where the spin-(1/2) is coupled to a macroscopic spin-(N/2) particle is also discussed.

Introduction. – Quantum heat engines (QHEs) are energy harvesting machines with quantum working substances. They allow for the exploration of the relations between energy and quantum coherence or correlations [1–14]. A particularly intriguing class of QHEs can harvest quantum coherent resources to produce useful work in a single thermal environment without violating the second law [11]. A major obstacle against their realization is quantum decoherence [9]. The latter can be beaten by scaling up quantum coherence in multi-level atoms [8] or by using superradiant atomic clusters [15]. Such schemes, however, involve sophisticated coherence-injection schemes with high energy costs. Moreover, the coherence needs to be small for quasi-thermal equilibrium operation, which makes the efficiency of the engine low. While the “low coherence” condition can be relaxed [16], finding simple and efficient preparation schemes of sizeable coherences remains a challenge against an efficient QHE.

In this letter, we consider quantum coherences, produced by interactions in spin systems at thermal equilibrium, as a scalable “natural quantum fuel” for a photonic Carnot engine (PCE) [11]. Such coherences can manifest themselves as thermal entanglement [17] or quantum discord [18], as verified experimentally [19,20]. Their robustness against decoherence has been studied [21–23]. The energetic cost of quantum correlations, refined out of coherences, is paid naturally here, in contrast to “artificial quantum fuels” schemes based upon coherence injection [11].

Our spin system is a star network in thermal equilibrium where a center spin is coupled to a set of N outer ones (cf. fig. 1) with Heisenberg XXZ interaction [24,25]. The state of the central spin has no coherence and is described by an effective temperature depending on the quantum coherence built in the network and scaled with N. When the central spins are used to power up a PCE, the coherences built in the network are exchanged exclusively as heat with the cavity field. In contrast to coherence-injection schemes, network coherences can be large as the central spin is in an effective thermal state. An isotropic model with N = 2 has been discussed in ref. [3].
Spin-star networks find applications in many areas ranging from the study of quantum decoherence [26,27] and spin baths [28,29] to quantum communication and cloning [24,30,31]. A closed spin-star network is a non-Markovian system and even a thermal state of the outer spins cannot thermalize the central one [32,33]. Spin-star networks could be realized by NMR molecules [34,35], superconducting qubits [36,37] or coupled microcavities [28,29]. At variance with most of the existing literature, which focuses on the quantum coherence of the central spin, our purpose exploits the thermal character of the central spin reflecting the coherences in the network. We consider an open network in thermal equilibrium with a hohlraum at temperature $T_b$. In this case, the outer spins are not at a thermal state while the central spin can be in an effective thermal state, at a different effective temperature than $T_b$. For two spins at low (high) temperatures, such temperature discrepancy is attributed to quantum entanglement (discord) [3].

We also consider an alternative system where a single spin-$(1/2)$ is coupled to a spin-$S$ particle ($S = 1/2, 1, 3/2, \ldots$) by the Heisenberg XXZ interaction [7], which can be compared to the case of a spin-tape coupled to a quantum-dot spin valve [39] (cf. fig. 1). The existence of thermal entanglement in this model, which could be implemented in molecular nanomagnets [40], superconducting qudits [41], or NV centers [42,43], has been addressed recently [44].

**Initialization.** – The two different spin environments that we consider to initialize the central spin-(1/2) particles before injecting them into a micromaser cavity [45] are depicted in fig. 1. The top of the right-most part of fig. 1 describes the scheme where spin-(1/2) particles in a spin-tape interact sequentially with a spin-$S$ particle at rest. In the other scheme, depicted at the bottom of the right-most panel of fig. 1, spin-tape particles interact sequentially with a spin-$S$ particle before injecting them into a micromaser cavity [45]. The magnitude $S$ of spin-$S$ and the number of the outer spins $N$ in the central spin network are related to each other via $N = 2S$. The parameters used for the plot are the Bohr frequency of the spins $\omega = 6.0$, spin-spin interaction coefficient $J = 0.8$ and anisotropy ratio $\lambda = 0.75$. All the parameters are dimensionless. $\omega$, $T_q$, and $J$ are scaled with $T_b$ (we have taken $\hbar = k_B = 1$). Atomic decay $\gamma$ is ignored in the plot.

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The operators of the stationary spin are denoted by $S_{1\alpha}$. Identifying $S_{1\alpha}$ as the collective spin operators $S_{1\alpha} = (1/2)\sum_{\alpha} S_{1\alpha}$, the two models can be intuitively considered as equivalent to each other. Even the dimensions of the two Hamiltonians are not the same, the thermodynamical interpretation of the local temperatures of the tape spins for both models are the same at $S = N/2$ and vary nonlinearly with $S$ as shown in fig. 1.

The density matrix of the total system for both cases is the thermal state $\rho = \exp (-\beta_b H)/Z$, where $Z = \text{Tr} [\exp(-\beta_b H)]$ is the partition function and $\beta_b = 1/T_b$ is the inverse bath temperature ($k_B = 1$). The reduced
density matrix $\rho_0$ of a tape spin is obtained by tracing out the other spin degrees of freedom and for both cases it is given by $\rho_0 = p_e |e\rangle \langle e| + (1 - p_e) |g\rangle \langle g|$, where $|e\rangle$ ($|g\rangle$) is the excited (ground) state of the spin-(1/2) and $p_e$ ($p_g = 1 - p_e$) is the corresponding occupation probability. For $p_e < 1/2$, a finite positive effective temperature for the tape spin can be defined as $T_q = -\omega/\ln(p_e/p_g)$ with $p_e = \exp(-\omega/2T_q)/Z_0$ and $Z_0 = 2 \cosh(\omega/2T_q)$. For the non-interacting case ($J = 0$), we always have $T_q = T_b$. However, $T_q$ can be different from $T_b$ in the presence of interactions, which can be used as an artificial hot bath to harvest work in a single thermal environment. For an interacting case, $T_q$ exhibits a nonlinear scaling with $N$ as shown in the bottom left plot in fig. 1. This effect can be exploited to power up a PCE.

**Photonic Carnot engine.** – A PCE [11] consists of a single-mode micromaser cavity acting as a working substance and undergoing a quantum Carnot cycle consisting of two quantum adiabatic and two isothermal processes [10]. The pump atoms injected into the cavity serve as an effective hot reservoir, while the environment of the system provides a cold bath. Radiation pressure of the cavity field performs work on one of the cavity mirrors playing the role of a mechanical piston. In contrast to the original proposal of PCE in which the pump atoms are in quasi-thermal states with weak coherence [11], we consider pump atoms in effective thermal states without any coherence. As seen above, the effective temperature $T_q$ of the atoms, however, depends on the number of spins in the ensemble and the interaction parameters. The discrepancy between such temperature and $T_b$ can be explained by the contributions of thermal quantum correlations [3].

During the isothermal expansion, the pump spins are sent through the cavity; the field is kept in a thermal state at temperature $T_q$ by changing the cavity field frequency, thus realising an effective expansion. For $p_e < 1/2$, the field in a micromaser is exactly thermal [46]. Assuming the cold bath is at the same temperature $T_b$ as the hohlraum, the thermal efficiency of the PCE can be given as the generalized Carnot efficiency $\eta = 1 - T_b/T_q$.

We determine the PCE dynamics by using two different approaches. First, we use a coarse-grained master equation, which allows for analytical results for ideal conditions. The second approach is numerical and entails the solution of the microscopic master equation, including the atomic decay and cavity loss. Initially, the cavity field is assumed to be in thermal state at $T_b$. The interaction with a tape spin is given by the Jaynes-Cummings Hamiltonian [47] $\hat{H}_I = g (\sigma_+ a + \sigma_- a^\dagger)$, where $g$ is the coupling constant, $\sigma_\pm = (\sigma_x \pm i \sigma_y)/2$ are the spin-(1/2) ladder operators for the tape spin, $a$ and $a^\dagger$ are the ladder operators of the cavity field with the unperturbed Hamiltonian $\hat{H}_c = \Omega a^\dagger a$, where $\Omega$ is the cavity field frequency. We consider a resonant interaction $\omega = \Omega$. Similar to the cycles in refs. [8,9,11], we assume $\Omega$ and $\omega$ change slightly during the isothermal stages.

For a large number of pump atoms, the atomic beam can be described as an effective hot bath using a coarse-grained master equation approach. If the atoms are injected randomly with arrival probability $r$ and an atom interacts with the cavity for a time interval $\tau (r \leq 1/\tau)$, then the coarse-grained master equation describing the evolution of the field reads [3]

$$\dot{\rho}_e = \frac{p_e'}{2} (2a^\dagger a \rho_e - \rho_e a^\dagger a^\dagger) + \frac{p'_g}{2} (2a^\dagger a^\dagger \rho_e - \rho_e a^\dagger a) , \tag{3}$$

where $\alpha = r (\gamma r)^2$, $\rho_e = p_e \rho_e$, $\rho'_e = p_g \rho_e$. Here $\rho_e$ and $\rho'_e$ are the effective rates for the excitation and deexcitation of the field, respectively. Below maser threshold, which is consistent with the detailed balance condition, the steady-state solution of eq. (3) is a thermal state at temperature $T_f = T_q = -\omega/\ln[p_e/(1-p_e)]$. Figure 2 shows how $T_f$ depends on $S = N/2$ for different $\lambda$. A nonlinear behaviour with $S$ emerges and it is tunable with $\lambda$ as tabulated in table 1. We observe deviation from linear scaling after $\lambda = 0.5$ in fig. 2(a). For $\lambda < 0.5$, linear scaling is the dominant behavior (cf. fig. 2(b)).

The nonlinear enhancement of $T_q$ with $N$ in fig. 2 can be related to quantum correlations. Though there are results of $N$-dependence of zero-temperature entanglement in XX model ($\lambda = 0$) [48], there are no proper measures of quantum correlations such as thermal discord and entanglement for multiple spin mixed states. Besides, while
Table 1: The fitting parameters to the data in fig. 2. A polynomial $y = a + bx + cx^2 + dx^3 + ex^4$ is used as a fitting curve for different anisotropy parameters $\lambda$. The coefficient of determination $R^2$ values are shown in the last column.

| $\lambda$ | $a$     | $b$     | $c$     | $d$     | $e$     | $R^2$ |
|-----------|---------|---------|---------|---------|---------|-------|
| 0.0       | 1.0030  | 0.0236  | 0.0     | 0.0     | 0.0     | 0.9995 |
| 0.25      | 0.9856  | 0.0708  | 0.0     | 0.0     | 0.0     | 0.9975 |
| 0.5       | 1.0215  | 0.0637  | 0.0138  | 0.0     | 0.0     | 0.9995 |
| 0.75      | 0.9552  | 0.2212  | -0.0394 | 0.0096  | 0.0     | 0.9997 |
| 0.9       | 0.8189  | 0.5149  | -0.1719 | 0.0306  | 0.0     | 0.9982 |
| 1.0       | 1.3518  | -0.7877 | 0.7572  | -0.2138 | 0.0222  | 0.9989 |

Fig. 3: (Colour online) Dependence of the effective temperature $T_q$ of the qubit on the cold bath temperature $T_b$. The parameters are the same as those used in fig. 1. The upper curves are for a larger number of outer spins $N$.

bipartite discord and concurrence contribute to enhancement of $T_f$ for the case of two spins [3], multipartite correlations mediated by the central spin [48] could contribute as well in the case of a large spin bath. Assuming that the correlations mediated by the central spin are relatively weak, we can focus on the pairwise coherence and correlations between the central spin and one of the outer spins. The behavior of $T_q$ against $T_b$, for different choices of $N$ is shown in fig. 3. Other parameters are as in fig. 1 and fig. 2. For $T_b \gtrsim \omega$, $T_q$ changes linearly with both $T_b$ and $N$ and, for $T_b \gg \omega$, we get $T_q = T_b$ for any $N$. The nonlinearity of $T_q$ with $T_b$ and $N$ is prominent for $T_b \sim 1$–4. We have calculated the concurrence in a central–outer spin pair and found that there is no entanglement in this regime. Only at $T_b \sim 0.5$, thermal entanglement is found. Remarkably, a similar behavior is found for the $C_1$-coherence in the state of the pair being considered (cf. fig. 4). Here, $C_1$ is the $l_1$-norm coherence [49], measuring the distance of the reduced density matrix of the central spin and one of the outer spins from their corresponding incoherent state. $C_1$ is independent of $N$ at high temperatures and decays slowly to zero after $T_b \sim 5$. In a regime of intermediate temperatures, where the nonlinear nature of $T_q$ against $N$ is prominent, $C_1$ appears to be also nonlinear with the number of outer particles. In line with $T_q$, $C_1$ increases linearly with the number of outer particles.

Fig. 4: (Colour online) Dependence of the concurrence $C_1$ (upper edges of panels) on the cold bath temperature $T_b = 1$ (front panel), 1, 2, 2.5, 3, 4, 6, 8, 10 and on the number of outer spins $N = 2, 3, 4, 5, 6$. The parameters are the same as those used in fig. 1.

The lowest-order polynomial that can fit well ($R^2 \sim 1$) the observed behaviour of $C_1$ at $T_b = 1$ is quadratic. We can find entanglement at higher temperatures by considering larger $J$. At $J = 1.6$ thermal entanglement emerges around $T_b = 1$, as shown in fig. 5, nonlinear with respect to $N$. However, in this case the nonlinear behavior of $T_q$
Spin-interaction–based work extraction

Fig. 6: (Colour online) Dependence of the effective temperature $T_q$ of the qubit on the cold bath temperature $T_b$. The parameters are the same as those used in fig. 1 except for $J = 1.6$. The upper curves are for a larger number of outer spins $N$.

Fig. 7: (Colour online) Dependence of the $l_1$-norm coherence $C_1$ (upper edges of panels) on the cold bath temperature $T_b$ at $N = 1$ (front panel), 1.5, 2, 2.5, 3, 4, 6, 8, 10 and on the number of outer spins $N = 2, 3, 4, 5, 6$. The parameters are the same as those used in the plot in fig. 1 except for $J = 1.6$.

with $T_b$ and $N$ is enhanced as shown in fig. 6 and similar to that of $C_1$ shown in fig. 7 (the lowest-order polynomial that can be used to fit $C_1$ at $T_b = 1$ is cubic). Our results supports the view that coherences are the main quantum resource for a PCE [4], but also suggest that correlations could still be significant as a more refined quantum resource. A spin-star quantum fuel is capable of “remote superradiance”, mediated by central spin, with more advantageous scaling with the number of quantum resources than superradiant [15] or multilevel quantum fuels [8]. The latter lead to $\sim N^2$ enhancement of $T_f$, while the spin-star allows for higher superlinear behavior.

Our PCE can operate with a single genuine heat bath as the other PCEs [8,11], without violating thermodynamical laws, as there is still an effective second bath generated by the central spin beam, which is in principle a non-equilibrium system. Effective temperature of the central spin is enhanced by thermal coherences whose costs are paid by the genuine heat bath. The efficiency of the spin-star scheme is several orders of magnitude larger than other PCEs [8,11]. For instance, for $\lambda = 1$ and $S = 5$, the cavity field temperature is $T_f / T_b \approx 3.50$ corresponding to $\eta \sim 70\%$. This estimation, however, is for ideal conditions, where we ignored the cooling of the cavity by the environment when it is empty during the initialization of the tape spins. Even if we inject the tape spins into the cavity immediately after they are ejected out of the hohlraum, the thermalization of the spin system would occur in finite time in real systems. In order to estimate how much time one would have for thermalisation without losing benefits of the scaling up the cavity temperature, we resort to a numerical analysis.

We apply the microscopic master equation approach in micromaser dynamics [50, 51], using typical microwave cavity parameters [52]. Instead of random injection of the atomic beam, we consider a regular atomic pump, i.e., atoms entering the cavity at equidistant times and an injection rate $r$. In some regimes, pump statistics can influence the variance of photon distribution [53]. Here we focus on the mean number of cavity photons, which determines the radiation pressure and the work output. The microscopic master equation approach simulates two stages of the micromaser dynamics. In the first one, atoms pass through the cavity one by one; each atom interacts with the cavity for a short time $\tau$. In the second stage, the cavity is empty for duration $\tau_0$ during which the tape spins are initialized and transferred to the cavity. Therefore, $1/r = \tau + \tau_0$. In order to achieve a steady state we have to “kick” the cavity field many times by repeated atomic injections before the field decays. We thus consider the number $N_{ex}$ of atoms interacting with the cavity field within the photon lifetime. Clearly, $1/r = 1/(N_{ex}\tau)$.

The initial density matrix of the composite system is given by the tensor product $\rho = \rho_c(0) \otimes \rho_0$, where $\rho_0$ is the effective atomic thermal state at $T_0$, initialized by either of the two schemes, and $\rho_c(0)$ is the thermal state of the cavity field at $T_c$. We take the cavity initial state at $T_c$ instead of $T_q$ so as to examine the effective quantum heating of the cavity by the tape spins. This allows us to see the capabilities of the spin-star fuel both for quantum thermalization [54] and for maintaining the cavity at a steady temperature. The initial state is evolved for the short time $\tau$ by the master equation [50, 51]

$$\dot{\rho} = -i[H, \rho] + \gamma D[\sigma_{0-}] + \kappa D[a],$$

where $H = H_c + (\omega/2)\sigma_0 + H_f$ is the total Hamiltonian of the system and $D[x] = (2\rho x^\dagger - \{x^\dagger x, \rho\})/2$ is the Liouvilian superoperator describing the relaxations to environment reservoir. The second term in eq. (4) is due to the atomic transition occurring without emitting a photon into the cavity mode at a rate $\gamma \approx 2\pi \sim 33.3$ Hz. We neglect the atomic dephasing term $(\gamma/2\pi \sim 3.3$ Hz). After that, the reduced density matrix of the cavity field is evolved by the master equation for the time $\tau_0$ [50, 51]

$$\dot{\rho}_c = -i[H_c, \rho_c] + \kappa D[a].$$
which characterizes the damped oscillatory behavior of the empty cavity between successive tape spins. The procedure is repeated by resetting the atomic state and updating the cavity field in the initial state, until the steady state of the cavity is reached. The master equations are written under the assumption of $T_b \ll \Omega$. In our simulations we take $\omega/T_b \sim 6$. For $\omega/2\pi = 50$ GHz this gives $T_b \sim 300$ mK. The excitation number in the cold bath $(n_b \sim 0.002)$ can be neglected and the cold bath treated effectively as a zero-temperature bath. For simplicity, we use the conditions in refs. [8,9,11], where the variation of $\Omega$ in isothermal stages is much smaller than $\Omega$ so that the parameters $\omega, \Omega, \gamma, \kappa$ can be considered constant in successive iterations of eqs. (4) and (5).

We calculate the steady-state value of the mean photon number in the cavity $\bar{n}_c = \text{Tr}(\rho_c a^\dagger a)$ to determine $T_f = \Omega / \ln(1 + 1/\bar{n}_c)$. The field temperature $T_f$ is effectively defined in the presence of cavity losses [54]. It gets closer to a genuine temperature when $\tau \ll 1/\omega$ [54]. We have $g/\pi = 50$ kHz which leads to $\tau \approx 40 \mu$s. We fix $\tau = 9.5 \mu$s and vary $\tau_0$ through $N_{\text{ex}}$. In fig. 8 we present the dynamics of $T_f$ depending on different $N_{\text{ex}}$ values. Initially the cavity is in thermal equilibrium with the cold bath $T_f/T_b = 1$. After injection of many spins at $T_q > T_b$, $T_f$ increases and reaches a steady-state value, which is smaller than $T_q$ due to losses. The coarse-grained master equation predicts a thermal state of the cavity at $T_f = T_q$. By increasing $N_{\text{ex}}$, the cavity interacts with more spins in a photon lifetime. Accordingly, the loss can be compensated. Around $N_{\text{ex}} \sim 6500$ the numerical behavior of $T_f$ conforms to the analytical predictions. In this case, time between successive spins becomes $1/\tau \sim 10 \mu$s for $\kappa \sim 5 \pi$ Hz. The cavity loss $\kappa = \omega/Q$ is evaluated for a high-finesse cavity with quality factor $Q = 2 \times 10^{10}$ at a microwave cavity mode frequency $\omega/2\pi = 50$ GHz. As we fixed the transition time at $\tau = 9.5 \mu$s, this leaves about $\tau_0 \sim 500$ ms for thermalization of the spin-star network and to take away the central spin to the cavity. If this time is found to be too short in a particular implementation, one can use an ensemble of spin-star networks, instead of a single one, and collect the central spins on demand to generate a convenient injection time distribution.

The harmful effect of increasing the empty cavity time $\tau_0$, or equivalently decreasing $N_{\text{ex}}$ can be studied by considering larger relaxation parameters. By fixing $\gamma$ and increasing $\kappa$, the advantages of scaling in $T_f$ with $S = N/2$ diminish, as shown in fig. 9(a). Similar conclusion can be deduced for increasing $\gamma$ at fixed $\kappa$ in fig. 9(b); besides, the sensitivities of the system to either atomic or cavity losses are similar, as revealed by comparing the top and bottom panels in fig. 9. Robustness of nonlinearity under strong damping can be exploited to tolerate longer thermalization and transfer times. We conclude that spin-star network-powered PCE can attain high efficiencies, even in the presence of decoherence. Note, however, that, even though a spin-star network can be considered as a natural quantum fuel, there could be other hidden costs to reduce the operational efficiency, such as coupling or decoupling the central qubit to the outer spins, which might require further optimizations, depending on specific implementation settings [55].

**Conclusions.** – We proposed the use of an $N$-element spin-star network as a quantum fuel to power up a PCE.
Such a choice offers significant advantages embodied by substantial quantum coherences generated and preserved naturally, in the form of thermal entanglement. When the network is in thermal equilibrium with a heat bath, the outer spins are in a state of non-equilibrium, and the central spin is in an effective thermal state with a temperature ten times greater than the one of the bath. Owing to the thermal nature of the central spin, large quantum coherences can be exploited to power up a PCE operating in contact with a single heat bath, even in the presence of atomic and cavity damping. We have found an efficiency of \( \sim 70\% \), which is several orders of magnitude larger than typical PCEs. In addition, we have discussed an equivalent scheme, where the outer spins are replaced by a single macroscopic spin. Various physical platforms, from quantum dots to molecular nanomagnets, superconducting spins, NMR and circuit-QED, can be considered for a proof-of-principle realization.

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ÖEM gratefully acknowledges support by Koç University (KU) Visiting Scholar Program by KU Office of VPAA (Vice President of Academic Affairs) and hospitality at Queen’s University Belfast. MP is supported by the EU FP7 grant TherMiQ (Grant Agreement 618074), and the J. Schwinger Foundation (grant No. JSF-14-7-0000). DT, FA and ÖEM acknowledge support from the University Research Agreement between Koç University and Lockheed Martin Corporation. MP and ÖEM acknowledge support from the Royal Society-Newton Mobility Grant N160057.

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