Thermodynamic cycle in a cavity optomechanical system

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Abstract
A cavity optomechanical system is initiated by the radiation pressure of a cavity field onto a mirror element acting as a quantum resonator. This radiation pressure can control the thermodynamic character of the mirror to some extent, such as by cooling its effective temperature. Here, we show that by properly engineering the spectral density of a thermal heat bath that interacts with a quantum system, the evolution of the quantum system can be effectively turned on and off. Inside a cavity optomechanical system, when the heat bath is realized by a multi-mode oscillator modelling of the mirror, this on–off effect translates to infusion or extraction of heat energy in and out of the cavity field, facilitating a four-stroke thermodynamic cycle.

Keywords: cavity optomechanical systems, quantum control, quantum thermodynamics

(Some figures may appear in colour only in the online journal)

1. Introduction
The study of cavity electrodynamics roughly began with the discovery of the Fabry–Perot interferometer, in which two transflective side mirrors sandwich an optical cavity of fixed length, thereby trapping an optical cavity field of designated wavelength inside. When one of the side mirrors is allowed to oscillate, usually by depositing a transflective surface on a micro cantilever, the trapped cavity field will interact with the movable mirror through radiation pressure [1] and other effects induced by the variable cavity length [2].

These kinds of controllable interactions provide some degree of manipulation to the movable mirror, thus opening the field of cavity optomechanics. In particular, extensive studies have been conducted during the last decade on how to cool down the effective temperature of the mirror [3–7]. Recently, studies on cavity optomechanical systems have found a wide range of applications such as quadrature squeezing of polaritons [8], generation of Kerr nonlinearity [9] and distant state entanglement [10, 11]. When the optomechanical coupling is so strong that photon blockades arise in the mirror, nonclassical and entangled states such as NOON and Bell states can be created [12, 13]. However, the quantum thermodynamic aspect of a cavity optomechanical system is less understood. In this article, we study the thermodynamic evolution of the cavity field under the influence of radiation pressure fed back from the mirror.

Generally speaking, when a heat bath modelled on an abstract manifold is coupled to a quantum multi-level system, thermodynamic adiabatic processes can be identified, during which work can flow in and out of the heat bath [14]. Further, if the manifold is assumed to be made up of spin systems with particular temperature gradients, thermodynamic machines can be facilitated [15, 16].

In the following sections, we show that these thermodynamic processes can be realized in a cavity optomechanical system, in which the mirror plays the role of a thermal controller that mediates heat energy in and out of the optical cavity. In fact, when modelling the mirror as a multi-mode quantum oscillator, we find the ensemble average energy of the cavity system dependent on the spectral density of states of the mirror. The mirror thermalizes the cavity system when its spectral density is specified such that the average energy evolves over time in the form of a square wave, giving off an on–off effect about the interaction between the cavity and the mirror. The jumping of the energy up and down on the square wave matches the endothermic and exothermic processes during which heat is either infused into or extracted.
from the system. The time during which the energy stays fixed
designates the adiabatic processes where no heat is transferred
but work is done on the cavity field, thus furnishing a full
thermodynamic cycle.

To understand this complex process more clearly and to
see why energy cycling can arise in the cavity, we study the
model in three steps, from simple toy models to finally a
realistic optomechanical system. First, we model the mirror
as a single-mode resonator with bosonic operators \( a \) and \( a^\dagger \)
and study how it would cycle energy in and out of a two-
level system \( \sigma \), in section 2. Since the usual interaction (either
radiation pressure or bolometric force) between a cavity field
and the reflective surface on a mechanical resonator does not
involve the absorption or radiation of photons, we model the
interaction by the Hamiltonian \( (a + a^\dagger)\sigma \).

We then extend the study to the scenario where the mirror
is modelled by a discrete set of resonator modes \( \{a_j\} \) in
section 3, while retaining the two-level system and the form of
its interaction to the mirror. We propose a method to construct
a square wave for the evolution of the energy of the two-level
system and test the viability of the construction through a
numerical simulation of the spectrum of the resonator modes.

Eventually, in section 4, we extend the resonator modes
for the mirror to a continuous spectrum and take the Dicke
limit for the number of two-level systems, mapping \( \sigma \) to
the bosonic number operator \( b^\dagger b \) to describe the cavity
field and transforming the toy model coupling to a realistic
radiation pressure \( b^\dagger b(a_j + a_j^\dagger) \). Modelling the heat bath on
the multi-mode mirror oscillator, the four processes of the
thermodynamic cycle are identified and the relevant spectral
density is given. The conclusion is given in section 5.

2. Single-mode oscillating effect

We consider a two-level system \( \sigma = |e\rangle \langle e| - |g\rangle \langle g| \) as
the main system and a single-mode oscillator \( \{a, a^\dagger\} \) as the
controller with a heat reservoir with the Hamiltonian \( \hbar = 1 \)
\[ H = \Omega \sigma + \omega a^\dagger a + \eta (a + a^\dagger) \sigma. \]  (1)

Let \( |\psi^+_n\rangle (|\psi^-_n\rangle) \) be the eigenstate of the controller associated
with the excited state \( |e\rangle \) (ground state \( |g\rangle \)) of the system, where \( n \)
designates the Fock number of the oscillator. The tensor product \( |e, \psi^+_n\rangle = |e\rangle \otimes |\psi^+_n\rangle \) describes the eigenstate
of the combined system and controller. Applying the Hamiltonian
(1) to this product state, we get
\[ H|e, \psi^+_n\rangle = [\Omega + \omega a^\dagger a + \eta (a + a^\dagger)]|e, \psi^+_n\rangle \]
\[ = \left[ \Omega + \omega \left(a^\dagger + \frac{\eta}{\omega}\right) a + \frac{\eta^2}{\omega} \right] |e, \psi^+_n\rangle. \]  (2)

This means that when the system stays in the excited state,
the part of the Hamiltonian that determines the evolution of
the controller is effectively a displaced oscillator
\[ H^e = \omega a^\dagger A_e - \frac{\eta^2}{\omega}, \]  (3)
where \( A_e = a + \eta/\omega \). In other words, while interacting with
the excited system, the eigenstate of the controller is a displaced
Fock state: \( |\psi^+_n\rangle = D(z)|n\rangle \), where \( D(z) = \exp\left[\frac{z}{2} (a^\dagger - a)\right] \) denotes the displacement operator. Associated with this
eigenstate, the eigenvalue for the effective Hamiltonian (3) is
then
\[ H^e |\psi^+_n\rangle = \epsilon_n |\psi^+_n\rangle = \left(\omega n - \frac{\eta^2}{\omega}\right) |\psi^+_n\rangle \]  (4)
and the total eigenenergy of the combined system and controller is
\[ E_n^e = \Omega + \omega n - \frac{\eta^2}{\omega}. \]  (5)

Following the same considerations, the system ground state \( |g\rangle \) is associated with the inversely displaced Fock state
\( |\psi^-_n\rangle = D(-\frac{\eta}{\omega})|n\rangle \) of the controller. The effective Hamiltonian
for the controller is
\[ H^g = \omega A_g^\dagger A_g - \frac{\eta^2}{\omega} \]  (6)
with \( A_g = a - \eta/\omega \), for which the total eigenenergy differs from
the excited state only by the sign of the system eigenenergy,
i.e.
\[ E_n^g = -\Omega + \omega n - \frac{\eta^2}{\omega}. \]  (7)

We can now consider the system dynamics for its coupling
to the controller thermal reservoir. Assume that the initial
state has the controller retaining a Fock number \( n \) and the
corresponding density matrix is in a thermal equilibrium with
B ernoulli distribution
\[ \rho(0) = P_g|g, \psi^+_0\rangle|g, \psi^+_0\rangle + P_e|e, \psi^+_0\rangle|e, \psi^+_0\rangle. \]  (8)

The Hamiltonian (1) drives the evolution of the system and the
controller separately according to what we discussed above, i.e.
\[ \rho(t) = e^{-iHt} \rho(0) e^{iHt} \]
\[ = P_g|e, \psi^+_n(t)\rangle|e, \psi^+_n(t)\rangle + P_e|g, \psi^+_n(t)\rangle|g, \psi^+_n(t)\rangle \]  (9)
where we have denoted \( |\psi^+_n(t)\rangle = e^{-iHt}|\psi^+_n(0)\rangle \) and
\( |\psi^-_n(t)\rangle = e^{-iHt}|\psi^-_n(0)\rangle \).

The energies stored in the system \( S \) and the controller \( C \)
vary with time according to the initial Bernoulli distribution and
the system parameters but their total energy remains static if
the density matrix starts off from the initial state given in
equation (8). Taking \( |e, |g\rangle \) as the basis of the system and
\( \{|\psi^+_n\rangle, |\psi^-_n\rangle\} \) as the basis of the controller, we can verify the constancy of the total energy, i.e.
\[ \langle H(t) \rangle = \text{tr}_{S+C} \rho(t) [\Omega \sigma_z + \omega a^\dagger a + \eta (a + a^\dagger) \sigma_z] \]
\[ = \Omega(P_e - P_g) + P_e\epsilon_n^e + P_g\epsilon_n^g. \]  (10)

However, since the controller acting as a heat reservoir has
constant influx or outflux of thermal energy to and from the
two-level system, the energy of the system and its interaction
with the controller will not remain constant. The average taken
over the reduced density matrix of the controller gives
\[ \langle \Omega \sigma_z + \eta (a + a^\dagger) \sigma_z \rangle_C = P_g|\psi^+_n(0)\rangle \Omega + \eta (a_n(t) + a_n^\dagger(t))|\psi^+_n\rangle \sigma_z \]
\[ + P_e|\psi^-_n(0)\rangle \Omega + \eta (a_n(t) + a_n^\dagger(t))|\psi^-_n\rangle \sigma_z, \]  (11)
where \( a_n(t) \) and \( a_n^\dagger(t) \) are the annihilation operators in the
Heisenberg picture.
of the controller following the evolutions of the excited state and ground state respectively. The Hamiltonians (3) and (6) can then be considered to be effectively driving the dynamics of the controller since \[ [H(t), a_i(t)]=[H^c(t), a_i(t)] \] and 

\[ [H(t), a_j(t)]=[H^c(t), a_j(t)]. \]

The time-dependent operators can be expressed explicitly by their Heisenberg equations with respect to these Hamiltonians:

\[ \dot{a}_i(t) = -i\alpha a_i(t) -\eta, \]

\[ \dot{a}_j(t) = -i\alpha a_j(t) + \eta, \]

for which the system–controller interaction \(\eta\) is essentially a driving of the level populations towards opposite directions for the two system levels. We will see in the next section that this driving translates to energy transfers in and out of the system levels. Substituting the formal solutions to equations (14)–(15) into (11), we find

\[ \langle \sigma_\alpha + \eta (a + a^\dagger) \rangle = \Omega \sigma_\alpha + \sum_{\gamma \in \{e, g\}} P_{\gamma} \eta \langle \psi_{\gamma}^e \rangle |a(0)\rangle e^{-i\omega t} \]

\[ + a^\dagger(0) e^{i\omega t} + (-1)^{\gamma(2\eta/\omega)} |\psi_{\gamma}^g\rangle \]

where we let \(\gamma\) be either 1 to indicate the positive sign for the excited state or 0 to indicate the negative sign for the ground state. The ensemble average of the operator \(a^\dagger\) at the initial state can be written as a c-number with amplitude \(\alpha\) and phase \(\phi\)

\[ \langle a^\dagger(0) \rangle = \langle \psi_{\alpha}^e | \alpha^\dagger(0) | \psi_{\alpha}^g \rangle + \langle \psi_{\alpha}^g | \alpha^\dagger(0) | \psi_{\alpha}^e \rangle = \alpha \cos \phi. \]  

Therefore, the average energy \(\langle \sigma_\alpha + \eta (a + a^\dagger) \rangle\) becomes an oscillating value

\[ \left[ \Omega + 2\eta \alpha \cos (\omega t - \phi) + 2\eta^2 \left( P_e - P_g \right) \cos \omega t \right] \sigma_\alpha, \]

where the direction of the oscillation depends on the system state.

### 3. Multi-mode square wave on–off effect

We now extend the concepts introduced in the last section to a multi-mode oscillator. The Hamiltonian (1) becomes

\[ H = \Omega \sigma_\alpha + \sum_j \omega_j a_j^\dagger a_j + \sum_j \eta_j (a_j + a_j^\dagger) \sigma_\alpha, \]

where the system part \(H_S = \Omega \sigma_\alpha\) stays identical while the controller part \(H_C = \sum_j \omega_j a_j^\dagger a_j\) and the interaction part \(H_I = \sum_j \eta_j (a_j + a_j^\dagger) \sigma_\alpha\) extends to the summation over all modes indexed by \(j\). The associated eigenstate for the excited state becomes a tensor product

\[ |e, \{ \psi_{\alpha}^e \}\rangle = |e\rangle \prod_{\alpha} |\psi_{\alpha}^e \rangle \]

over the Fock states of all the photonic modes of the field.

Applying equation (19) to the eigenstate, we find the effective Hamiltonian for the multi-mode controller \(|\{ \psi_{\alpha}^e \}\rangle\)

\[ H^c = \sum_j \left( \omega_j A_{\alpha,j} A_{\alpha,j}^\dagger - \frac{\eta_j^2}{\omega_j} \right), \]

where the displaced annihilation operator is now distinguished for each mode

\[ A_{\alpha,j} = D^{-1} \left( \frac{\eta_j}{\omega_j} \right) a_j D \left( \frac{\eta_j}{\omega_j} \right). \]

An identical procedure can be applied to the ground state, for which the index \(e\) in the equations above will be replaced by \(g\).

When the two-level system retains its Bernoulli distribution, the associated density matrix here only differs from equation (9) by the expressions in the eigenstates, i.e. we can verify

\[ \rho(t) = P_e |e, \{ \psi_{\alpha}^e \}\rangle |e, \{ \psi_{\alpha}^e \}\rangle^\dagger + P_g |g, \{ \psi_{\alpha}^g \}\rangle |g, \{ \psi_{\alpha}^g \}\rangle^\dagger \]

where

\[ |e, \{ \psi_{\alpha}^e \}\rangle = |e\rangle \otimes e^{-i\omega t} |\{ \psi_{\alpha}^e \}\rangle \]

and similarly for the ground state. We can see that the partition of eigenstates here is no different from those given in equation (9) and therefore the evolution of ladder operators given in equations (14)–(15) can be naturally extended to the multi-mode case, where each \(a_j\) is now evolving through the displacements of all eigenmodes as described by equation (21) instead of each eigenmode alone.

That means that equations (14)–(15) still apply and the average energy over the orthogonal controller basis \(|\{ \psi_{\alpha}^e \}, |\{ \psi_{\alpha}^g \}\}\rangle\) is

\[ \langle H_S + H_I \rangle = \left[ \Omega + \sum_j 2\eta_j \left( \alpha_j \cos (\omega_j t - \phi_j) + \frac{\eta_j}{\omega_j} \left( P_e - P_g \right) \cos \omega_j t \right) \right] \sigma_\alpha, \]

which is just a multi-mode extension of the single-mode case given in equation (18). Similarly, \(\alpha_j\) and \(\phi_j\) are the initial amplitude and phase of the \(j\)th mode.

This extension, though simple, gives a large edge of control to the controller over the system. Note that, by reshuffling the terms, we can arrange it to become a Fourier series about time \(t\) with nonzero coefficients in both the sines and the cosines

\[ \langle H_S + H_I \rangle = \sigma_\alpha \left[ \Omega - \left( P_e - P_g \right) \sum_j \frac{2\eta_j^2}{\omega_j} \right. \]

\[ + \sum_j 2\eta_j \left( \alpha_j \cos \phi_j + \frac{\eta_j}{\omega_j} \left( P_e - P_g \right) \cos \omega_j t \right. \]

\[ \left. + \alpha_j \sin \phi_j \sin \omega_j t \right]. \]

When setting this Fourier series with different coefficients, we can obtain different cyclic waves. In other words, the multi-mode coupling between the oscillator as the controller and
the two-level system enables us to control the average energy stored in the system by setting different initial states for the controller.

The typical case is when setting the Fourier series as a square wave, for which the on–off switching of the stored energy occurs. Consider that we let

\[
\begin{align*}
2\eta_j\alpha_j \sin \phi_j &= \frac{A}{\alpha_j} \\
2\eta_j \left( \alpha_j \cos \phi_j + (P_e - P_g) \frac{\eta_j}{\alpha_j} \right) &= 0
\end{align*}
\]

(27)
to zero out the cosine series. Then the square wave of height \(A\) would be realized if \(\omega_j = (2j - 1)\omega_0\) are the odd harmonics of some fundamental frequency \(\omega_0\) and

\[
\alpha_j = \sqrt{\frac{A(P_e - P_g)^2\eta_j^4 + A^2}{2\eta_j(2j - 1)\omega_0}}.
\]

(28)

\[
\phi_j = -\tan^{-1} \left( \frac{A}{2(P_e - P_g)\eta_j} \right).
\]

(29)

We can observe and conclude that the amplitudes \(\alpha_j\) determine the frequency of the Fourier series while the phases \(\phi_j\) determine the height of the square wave or the amount of energy being transferred in and out of the system.

Finally, following the nomenclature of quantum thermodynamics, we can define a spectral density function

\[
J(\omega) = \sum_j \frac{4(P_e - P_g)^2\eta_j^4 + A^2}{4\eta_j^2\omega_j^2} \delta(\omega - \omega_j)
\]

(30)

for the oscillator controller as a heat reservoir. This system-specific reservoir will supply energy to the system such that it will undergo energy cycling with the system at period \(T = 2\pi/\omega_0\), during which for half of the time the system will be turned on to attain a higher energy and for half of the time the system will be turned off to a lower energy state. The distribution of the spectral density depends on the coupling strengths \(\eta_j\) as well as the desired population distribution \(\{P_e, P_g\}\) of the two-level systems.

In figure 1, the spectral density is plotted for three values of \(P_e - P_g\): -1 for no population inversion, 0.001 for half inversion (uniform Bernoulli distribution), and 1 for full population inversion. The targeted amplitude \(A\) is \(\frac{30}{2\pi}\) MHz. The base frequency \(\omega_0\) is set the experimentally accessible detuning of \(\frac{40}{2\pi}\) MHz between the cavity and a driving laser. The first 50 odd harmonics over this base frequency are considered for the heat reservoir. We assume the spectral density of the coupling strength to adopt a normal distribution with variance \(\frac{100}{2\pi}\) kHz about the mean \(\frac{10}{2\pi}\) MHz. We can note that to realize the on–off switching effect, the oscillator is generally not very difficult to prepare. Regardless of the amount of inversion the two-level system retains, the associated spectral density of the oscillator as a heat reservoir has fairly uniform phase distributions. As for the magnitude \(J(\omega)\), we observe that it is sufficient to produce only the first few harmonics for the purpose of heat transfer. In particular, for the half-inversion case, the amplitude of \(J(\omega)\) is almost negligible across all harmonics since the uniform Bernoulli distribution at the two-level system naturally induces the energy exchange with the oscillator controller.

4. Thermodynamic cycle of cavity optomechanical system

With the on–off effect shown on an oscillator-coupling two-level system, we now turn eventually to the study of a
movable mirror in an optomechanical cavity associated with the thermodynamic energy transfer. The previously studied two-level system $\sigma_z$ is replaced here by a pair of annihilation and creation operators $\{b, b^\dagger\}$ to represent an optical field traversed in the cavity. The effective dynamics of the movable mirror mounted on a cantilever is determined by the flexibility modulus of the cantilever materials. It is modelled by a multi-mode oscillator with Hamiltonian $\sum_j \omega_j a_j^\dagger a_j$ and can be regarded as a heat bath under the Feynman–Vernon [17] and Caldeira–Leggett [18] frameworks.

The motion of the mirror deforms the cavity volume, which results in a radiation pressure proportional to the cavity photon number $b^\dagger b$ and the mirror displacement $x$ being fed back to the mirror [11]. Expressing the displacement $x$ in terms of the canonical conjugate variables $a_j$ and $a_j^\dagger$, the total Hamiltonian reads

$$H = H_S + H_C + H_I = \Omega b^\dagger b + \sum_j \omega_j a_j^\dagger a_j + b^\dagger b \sum_j \eta_j(a_j + a_j^\dagger).$$

(31)

To examine the heat exchange process over time, we consider the density matrix

$$\rho(0) = \rho_S(0) \otimes \rho_C(0) = \sum_m P_m |m, \{\psi^m_n\}|m, \{\psi^m_n\}$$

(32)

to represent an initial mixed state at thermal equilibrium, where $|m\rangle$ is the Fock eigenstate for the photon number in the cavity and $|\psi^m_n\rangle$ is the associated eigenstate of the mirror controller. At time $t$, the effective Hamiltonian that drives the evolution of $|\psi^m_n\rangle$ is

$$H^m = m\Omega + \sum_j \omega_j a_j^\dagger a_j + m \sum_j \eta_j(a_j + a_j^\dagger),$$

(33)

which gives rise to the reduced density matrix of the mirror controller as

$$\rho_C(t) = \sum_m P_m e^{-iH^m t} |m, \{\psi^m_n\}|m, \{\psi^m_n\} e^{iH^m t}.$$  

(34)

Letting $\langle H_S(t) + H_I(t) \rangle_C = \Omega(t)b^\dagger b$, we find the effective eigenenergy of the cavity system to be

$$\Omega(t) = \Omega + \sum_m P_m \sum_j \eta_j |\psi^m_n\rangle [a_{m,j}(t) + a_{m,j}^\dagger(t)] |\psi^m_n\rangle.$$

(35)

where $a_{m,j}(t)$ and $a_{m,j}^\dagger(t)$ are the evolved operators in the Heisenberg picture similar to those defined in equations (12)–(13), except that the system states are here extended to all $m$ levels for each $j$th mode.

Following the same routine as in the previous section but assuming a continuous spectrum for the mirror as a heat reservoir, we would arrive at

$$\Omega(t) = \Omega_0 - 2 \int d\omega \eta_0 \frac{\eta_2}{\omega} \left(\alpha \cos \phi + \frac{\mu n}{\omega} \cos \omega t \right) \sin \omega t,$$

(36)

where $\Omega_0 = \Omega - 2\mu \int (\eta^2 / \omega) d\omega$ is the renormalized energy offset when the cavity interacts with the heat reservoir. In the equation, $\mu = \sum_m P_m m$ is the weighted mean of the photon population across all levels in the cavity system. To furnish the on–off effect studied in the last section, the phase and magnitude of the spectral density should therefore be

$$\phi(\omega) = -\tan^{-1} \frac{A \omega_0}{2 \mu \eta^2},$$

(37)

$$J(\omega) = \frac{4 \mu^2 \eta^4 + A^2 \omega_0^2}{4 \eta^4 \omega^2}.$$  

(38)

If the spectral density is set thus, $\Omega(t)$ will cycle about $\Omega_0$ with amplitude $A$ between two constant values, as illustrated in figure 2, where the reservoir is again assumed to consist of 50 odd harmonics over a base frequency. The durations within which $\Omega(t)$ stays constant can be identified with adiabatic processes. During these processes, though the cavity remains interacting with the reservoir, no energy is transferred into or out of the cavity and all levels $m$ remain constantly spaced.
Between these processes, the cavity system either absorbs energy from the reservoir, raising up all levels simultaneously, or infuses energy back to the reservoir, letting itself fall back to the original levels. The absorption of energy can be identified thermodynamically with an endothermic process, whereas the depletion of energy can be identified with an exothermic process.

5. Conclusion

We have shown a specific thermodynamic cycle on a cavity optomechanical system. By first demonstrating the cyclic eigenenergy of a two-level system interacting with a single-mode oscillator, we then prove that within the cavity system, the cavity field can act as the thermal system relative to the multi-mode oscillating mirror acting as the heat reservoir that controls the energy flow. By matching the spectral density of the mirror with that of a square wave in the cyclic eigenenergy, an on–off cycle of energy exchange can be constructed, during which four thermodynamic stroke processes can be identified.

Therefore, a quantum mechanical system with appropriate spectral densities can serve as a quantum thermodynamic machine. Future investigations will focus on how the constructed thermalization processes fit within the general framework of quantum heat engines.

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