Minimal temperature of quantum refrigerators

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Abstract – A first-principle reciprocating quantum refrigerator is investigated to determine the limitations of cooling to absolute zero. If the energy spectrum of the working medium possesses an uncontrollable gap, then there is a minimum achievable temperature above zero. Such a gap, combined with a negligible amount of noise, prevents adiabatic following during the expansion stage which is the necessary condition for reaching $T_c \to 0$.

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Introduction. – It is well accepted that cooling to the absolute zero becomes prohibitively difficult [1,2]. More precisely, is the absolute-zero state connected to quantum friction leadsto a finite minimal temperature. In the tradition of this issue. In the present study we want to show that there is quantum twist to this issue. In the tradition of thermodynamics we try to generalize from a specific example. Our previous analysis of quantum models have shown that the cooling rate vanishes as $T_\infty$ where $\delta > 1$ [3–7], i.e. $T_c = 0$ is reachable. The present study analyzes a reciprocating quantum refrigerator which is subject to noise on the external controls of the adiabatic expansion. We find that such noise is sufficient to make $T_c = 0$ unattainable.

Reciprocating refrigerators operate by a working medium shuttling heat from the cold to the hot reservoir. This requires external control of the temperature of the working medium. Upon contact with the cold side the working-medium temperature has to become lower than $T_\infty$ — the cold-bath temperature.

A generic working medium possesses a Hamiltonian that is only partially controlled externally:

$$\hat{H} = \hat{H}_{\text{int}} + \hat{H}_{\text{ext}}(\omega),$$

(1)

where $\omega = \omega(t)$ is the time-dependent external control field. Typically, $[\hat{H}_{\text{int}}, \hat{H}_{\text{ext}}] \neq 0$, therefore, $[\hat{H}(t), \hat{H}(t')] \neq 0$ as a result a state diagonal in the temporary energy eigenstates cannot follow adiabatically. This fact, which is the source of quantum friction, has a profound effect on the performance of the heat pump [8,9].

Almost perfect adiabaticity is the key to low-temperature refrigeration. Typically, the internal interaction leads to an uncontrollable finite gap $hJ$ in the energy level spectrum between the ground and first-excited state. We will show that this gap combined with unavoidable quantum friction leads to a finite minimal temperature.

The cycle of operation, the quantum heat pump. – The working medium in the present study is composed of an interacting spin system. Equation (1) is modeled by the $SU(2)$ algebra of operators. We can realize the model by a system of two coupled spins $\hat{H}_{\text{int}} = \frac{1}{2} h J (\hat{\sigma}_z \otimes \hat{\sigma}_z^2 - \hat{\sigma}_x \otimes \hat{\sigma}_x^2) \equiv hJ \hat{\mathbf{B}}_2$, where $\hat{\sigma}$ represents the spin-Pauli operators, and $J$ scales the strength of the interparticle interaction. For $J \to 0$, the system approaches a working medium with non-interacting atoms [5]. The external Hamiltonian represents the interaction of spins with an external magnetic field:

$$\hat{H}_{\text{ext}} = \frac{1}{2} h \omega(t)(\hat{\sigma}_z \otimes \hat{I}^2 + \hat{I}_1 \otimes \hat{\sigma}_z^2) \equiv h\omega(t) \hat{\mathbf{B}}_1.$$

The $SU(2)$ is closed with $\hat{\mathbf{B}}_3 = \frac{1}{2} (\hat{\sigma}_y \otimes \hat{\sigma}_x^2 + \hat{\sigma}_x \otimes \hat{\sigma}_y^2)$ and $\{\hat{\mathbf{B}}_1, \hat{\mathbf{B}}_2\} \equiv 2i \hat{\mathbf{B}}_3$.

The total Hamiltonian then becomes

$$\hat{H} = h \left( \omega(t) \hat{\mathbf{B}}_1 + J \hat{\mathbf{B}}_2 \right).$$

(2)

The temporary energy levels, the eigenvalues of $\hat{H}$ are $\epsilon_1 = -h\Omega$, $\epsilon_2/3 = 0$, $\epsilon_4 = h\Omega$, where $\Omega = \sqrt{\omega^2 + J^2}$. For $J \neq 0$ there is a zero-field splitting. Equation (2) contains the essential features of the Hamiltonian of magnetic materials [10].
adiabat and the compression
The hot and cold isotherms are indicated. Both the expansion
periods.

$S_{\text{VN}} - S_{p}$

\[ T \]

Fig. 1: (Colour on-line) Refrigerator cycle in the frequency
entropy plane. The von Neumann entropy
\[ S_{\text{VN}} = -\text{tr}\{ \rho \log \rho \} \]

\[ \text{ABCD} \]

in addition, there are two segments, termed
isochores, where the working medium is in contact
with the cold/hot baths and the external control field $\omega$ is
constant. In addition, there are two segments, termed
adiabats, where the external field $\omega(t)$ varies and with it
the energy level structure of the working medium. This
cycle is a quantum analogue of the Otto cycle [14]. Each
segment is characterized by a quantum propagator $U_c$.
The propagator maps the initial state of the working medium
to the final state on the relevant segment. The four strokes
of the cycle (see fig. 1) are

- **Isochore (isomagnetic)** $A \rightarrow B$: the field is maintained
  constant $\omega = \omega_h$, the working medium is in contact
  with the hot bath of temperature $T_h$. $L_D$ leads to equilibrium
  with heat conductance $G_G$ for a period $\tau_c$. The segment dynamics is described by
  the propagator $U_h$. The product of the four propagators,$U_c$ is the cycle
  propagator at point $A$:

\[ U_{\text{cyc}} = U_A U_B U_H U_C. \] (4)

Eventually, independent of initial condition, after a few
cycles, the working medium will reach a limit cycle character-
ized as an invariant eigenvector of $U_{\text{cyc}}$ with
eigenvalue $1$ (one) [9]. The characteristics of the refrigerator
are therefore extracted from the limit cycle.

**The dynamics of the expansion adiabat.** – The
key to low temperatures is the expansion adiabat. In
magnetic-salt–based refrigerators this segment is termed
the adiabatic demagnetization stage [10,15,16]. A neces-
sary condition for cooling is that the energy of the working
medium at contact point $C$ is lower than the equilibrium
energy at temperature $T_c$. What are the starting condi-
tions at the beginning of the expansion segment point $B$?
Considering that the efficiency is limited by the Carnot
cycle $\eta_{\text{otto}} \lessgtr \eta_{\text{carnot}}$ for a refrigerator this leads to the
reversed condition $\Omega_c \lessgtr \frac{T_c}{T_h}$. Now $\Omega_c \geq J$ therefore using
\[ \Omega_c(\min) = J: \]

\[ T_c \geq J \frac{T_h}{\Omega_h}. \] (5)

This condition relates the hot-end frequency $\omega_h$ to the
temperature $T_c$ [10]. To force $T_c$ to zero $\omega_h \rightarrow \infty$ and with it $\Omega_c$. Under these conditions at equilibrium all population
is in the ground state and $\langle \hat{H} \rangle_{h = 0} = -\hbar \Omega_c$. This is the
optimal starting point for the expansion adiabat. On the
cold side the necessary condition for refrigeration is that
the internal energy of the working medium at the end of the
expansion is smaller than the equilibrium energy with the
cold bath:

\[ \langle \hat{H} \rangle_{c} \lessgtr \langle \hat{H} \rangle_{\text{eq}}(T_c) = -\hbar \Omega_c \left( 1 - 2 e^{\frac{\hbar \Omega_c}{k_B T_c}} \right). \] (6)

where $\langle \hat{H} \rangle_{\text{eq}}(T_c)$ is approximated by the low-temperature
limit $\hbar \Omega_c \gg k_B T_c$. Such a condition is fulfilled if the
populations follows adiabatically the ground state during
the expansion adiabat. Then $\langle \hat{H} \rangle_c = -\hbar \Omega_c$ and eq. (6)
is fulfilled. The expansion stage requires to reduce the
external field $\omega$ from a large to a very small value
maintaining adiabaticity.

The orthogonal set of time-independent operators $\hat{B}_c$ is
closed to the dynamics, and therefore they can supply a
complete vector space to represent the propagators $U_{hc}$.
The adiabatic limit is described by $\mu \to 0$. Then eq. (10) factorizes. These are the perfect-adiabatic-following conditions. In general eq. (10) describes a periodic motion of $\hat{H}$, $\hat{L}$ and $\hat{C}$. Each period is defined by

$$q\Theta = 2\pi l \quad l = 0, 1, 2 \ldots , \quad (11)$$

where $l$ is the winding number. At the end of each period $U_{2l}$ restores to the identity matrix. These are the frictionless conditions of adiabatic following. For intermediate times $(\hat{H})$ is always larger than the frictionless value. The amplitude of this periodic dynamics decreases when $\mu$ becomes smaller, Cf. $U_{2l}(1, 1)$ in eq. (10).

The frictionless conditions define a quantization condition for the adiabatic parameter $\mu$:

$$\mu = \left(\frac{(2\pi l)^2}{\Phi_{hc}} - 1\right)^{-\frac{1}{2}} \quad (12)$$

Examining eq. (12) we find that there is no solution for $l = 0$. The first frictionless solution $l \geq \sqrt{\frac{2\pi\mu}{\Phi_{hc}}}$ leads to a minimum expansion time for frictionless solutions:

$$\tau_{hc}(\text{min}) = K_{hc} \sqrt{\frac{(2\pi l)^2}{\Phi_{hc}}} - 1 \quad (13)$$

The family of all frictionless solutions leads to refrigeration cycles which obey eq. (6) for any $T_c > 0$. Such frictionless refrigerators have no minimum temperature therefore they are connected to $T_c = 0$.

The effective minimal temperature. – Any realistic refrigerator is subject to noise on the external controls. Perfect adiabaticity requires precise control of the scheduling of the external field $\omega(t)$. Any deviation from perfect adiabatic following, maintaining the ground state on the expansion $\hat{H}$ will lead to a minimum temperature. If $\langle\hat{H}\rangle_c = -\Omega_\omega(1 - \delta)$, where $\delta$ is the deviation from perfect adiabatic following then, from eq. (6) $\delta \leq 2e^{-\frac{\hbar\Omega}{\Phi_{hc}}}$, leading to

$$T_c \geq \frac{\hbar\Omega_c}{-k_b \log(\delta/2)} \geq \frac{\hbar\Omega}{-k_b \log(\delta/2)} \quad (14)$$

From eq. (10), $\delta = \mu^2(1 - c)/(1 + \mu^2)$ which is zero for the periodic solutions. $\delta$ also vanishes when $\mu \to 0$. We will now show that even a significant amount of noise will lead to $\delta > 0$ and $T_c(\text{min}) > 0$.

First we consider a piecewise process controlling the scheduling of $\omega$ in time. At every time interval, $\omega$ is updated to its new value. Then random errors are expected in the duration of these time intervals described by the Liouville operator $L_N$. This process is mathematically equivalent to a dephasing process on the expansion adiabat [18]. This stochastic dynamics can be modeled by a Gaussian semigroup with the generator [13,19]

$$L_{N, s}(A) = -\frac{2p}{h^2} [\hat{H}, [\hat{H}, A]], \quad (15)$$

where $s = \sin(q\Theta)$ and $c = \cos(q\Theta)$. 

A more thermodynamically oriented alternative is based on a time-dependent set. The set includes the energy $\hat{H}$ and two other orthogonal operators:

$$\begin{align*}
\hat{H} &= \omega(t)\hat{B}_1 + J\hat{B}_2, \\
\hat{L} &= -i\hat{B}_1 + \omega(t)\hat{B}_2, \\
\hat{C} &= \Omega(t)\hat{B}_3.
\end{align*} \quad (7)$$

In general the dynamics on the expansion adiabat is generated by $L = L_H + L_N$, where $L_H = \frac{d}{dt}[\hat{H}, -] \text{ and } \hat{H}(t)$ the time-dependent Hamiltonian, eq. (2). The external noise generator is $L_N$ defined later. For perfect adiabatic following the propagator $U$ factorizes between $\hat{H}$, and the pair $\hat{L}$, $\hat{C}$.

The noiseless dynamics generated only by the Hamiltonian $\hat{H}(t)$ is the key to adiabaticity:

$$\frac{d}{dt} \left( \begin{array}{c} \hat{H} \\ \hat{L} \\ \hat{C} \end{array} \right) (t) = \left( \begin{array}{ccc} \Omega & -\frac{J\omega}{\Omega} & 0 \\ \frac{2\omega}{\Omega} & \Omega & -1 \\ 0 & 1 & \Omega \end{array} \right) \left( \begin{array}{c} \hat{H} \\ \hat{L} \\ \hat{C} \end{array} \right) . \quad (8)$$

The ability of the working medium to follow the energy spectrum is defined by the adiabatic measure $\mu = \frac{\omega(t)}{\Omega(t)}$. If $\mu = 0$ the propagator factorizes. Constant $\mu$ minimizes the non-adiabatic deviations during the expansion [17]. In addition constant $\mu$ leads to a closed-form solution for the propagator $U_{hc}$ forcing a particular scheduling of the external field $\omega(t)$ with time: $\omega(t) = Jf/\sqrt{1-f^2}$, where $f$ is a linear function of time: $f = \frac{1}{\tau_{hc}}(\frac{\Omega}{\Omega} - \frac{\omega(t)}{\Omega}) + \frac{\omega_0}{\Omega}$. The adiabatic parameter $\mu$ and the time allocated to the adiabat $\tau_{hc}$ obey the reciprocal relation $\mu = K_{hc} \tau_{hc}$, where $K_{hc} = \frac{1}{2}(\frac{\omega_0}{\Omega} - \frac{\omega(t)}{\Omega})$.

Equation (8) is integrated by defining a new time variable: $d\hat{t} = \Omega dt$. The final values of $\Theta_{hc}$ becomes $\Theta_{hc} = \tau_{hc} \frac{K_{hc}}{\Phi_{hc}} \Phi_{hc}$, where $\Phi_{hc} = (\arcsin(\frac{\omega(t)}{\Omega}) - \arcsin(\frac{\omega_0}{\Omega}))$ and $0 > \Phi > -\frac{\pi}{2}$.

Equation (8) is solved by noticing that the diagonal is a unit matrix multiplied by a time-dependent scalar. Therefore we seek a solution of the type $U_{hc} = \hat{U}_1 \hat{U}_2$, where $[\hat{U}_1, \hat{U}_2] = 0$. The integral of the diagonal part of eq. (8) becomes

$$\hat{U}_1 = e^{\int_0^{t_{hc}} \hat{H} dt} = \frac{\Omega_\omega}{\Omega_b} \frac{\Omega}{\Omega_b} I, \quad (9)$$

which can be interpreted as the scaling of the energy levels with the variation in $\Omega$.

To integrate $\hat{U}_2$, the non-diagonal parts of eq. (8) are diagonalized, leading to the eigenvalues $0, -i\sqrt{q}, i\sqrt{q}$, where $q = \sqrt{1 + \mu^2}$, and the propagator

$$\hat{U}_2 = \begin{pmatrix} 1 + q^2c & \frac{\mu^2}{q} & \frac{\mu(1-c)}{q^2} \\ \frac{\mu^2}{q} & c & \frac{-q}{q^2} \\ \frac{\mu(1-c)}{q^2} & \frac{q}{q^2} & q^2 + c \end{pmatrix} , \quad (10)$$

where $s = \sin(q\Theta)$ and $c = \cos(q\Theta)$.
which is termed phase noise. The modified equations of motion on the adiabats become

$$\frac{d}{dt} \begin{pmatrix} \mathbf{H} \\ \mathbf{L} \end{pmatrix} (t) = \begin{pmatrix} \frac{\Omega}{\hbar} & -\frac{\gamma p}{\hbar} & 0 \\ -\frac{\gamma p}{\hbar} & \frac{\Omega}{\hbar} & \gamma p \Omega & -1 \\ 0 & 1 & \frac{\Omega}{\hbar} - \gamma p \Omega \end{pmatrix} \begin{pmatrix} \mathbf{H} \\ \mathbf{L} \end{pmatrix}.$$  

(16)

We seek a product form solution: \(U_{hc} = U_1 U_2 U_3\), where \(U_1\) the noise propagator, given by \(\frac{d}{dt} U_3(t) = \mathcal{W}(t) U_3(t)\), where

\[
\mathcal{W}(t) = U_2(-t) \begin{pmatrix} 0 & 0 & 0 & \mu s & -\mu (1-c) \\ 0 & -\gamma p \Omega & 0 & \mu s & 1 & 0 \\ -\mu (1-c) & 0 & 1 & 0 \\ \end{pmatrix} U_2(t). \tag{17}
\]

We seek an approximate solution for \(U_3\) in the limit when \(\mu \to 0\), then \(U_2 = I\) since this is the frictionless limit. Expanding eq. (17) to first order in \(\mu\) leads to

\[
\mathcal{W}(t) \approx -\gamma p \Omega(t) \begin{pmatrix} 0 & 0 & 0 & \mu s & -\mu (1-c) \\ 0 & -\gamma p \Omega & 0 & \mu s & 1 & 0 \\ -\mu (1-c) & 0 & 1 & 0 \\ \end{pmatrix}. \tag{18}
\]

\(U_3(\gamma_{hc})\) is solved in two steps. First evaluating the propagator for one period of \(\Theta\), for which \(\Omega(t)\) is almost constant, and then the global propagator becomes the product of the one-period propagators for \(l\) periods: \(U_3(\gamma_{hc}) \approx U_3(\Theta = 2\pi)^l\). The Magnus expansion to second order is employed to obtain the one-period propagator \(U_3(2\pi)\):

\[
U_3(\Theta = 2\pi) \approx e^{M_1 + M_2 + \ldots}, \tag{19}
\]

where \(M_1 = \int_0^{2\pi} d\Theta \mathcal{W}(\Theta)\) and \(M_2 = \frac{1}{2} \int_0^{2\pi} \int_0^{2\pi} d\Theta d\Theta' [\mathcal{W}(\Theta), \mathcal{W}(\Theta')] + \ldots\). The first-order Magnus term leads to

\[
U_3(\Theta = 2\pi) M_1 \approx \begin{pmatrix} 0 & 0 & \mu (1 - e^{-2\pi \gamma p \Omega}) \\ 0 & 1 & 0 \\ e^{-2\pi \gamma p \Omega} & 0 & e^{-2\pi \gamma p \Omega} \\ \end{pmatrix}, \tag{20}
\]

for which to first order in \(\mu, \delta\), the deviation from perfect adiabatic following, is zero. The second-order Magnus approximation leads to

\[
U_3(\Theta = 2\pi) M_2 \approx \begin{pmatrix} C & -S & 0 \\ S & C & 0 \\ 0 & 0 & 1 \\ \end{pmatrix}, \tag{21}
\]

where \(S = \sin \alpha\) and \(C = \cos \alpha\). \(\alpha = \pi \gamma p \Omega \sqrt{y^2 + 4}\) and as \(\mu \to 0\), \(\alpha = 2\pi \gamma p \Omega \mu\). The second-order propagator \(U_3(\gamma_{hc})\), for \(l\) revolutions is also a rotation matrix identical to eq. (21), with a new angle \(\alpha_l = 2\pi \gamma p \mu \int_0^{2\pi l} \Omega(\Theta) d\Theta = \pi \gamma p J \ln \left[ \frac{\Omega (\Theta + \omega_l)}{\Omega (\Theta - \omega_l)} \right] \frac{v}{2l}. \]

The deviation of \(U_3\) from the identity operator defines \(\delta\). Asymptotically for \(\mu \to 0\), \(\delta_{min} = 1 - \cos(\alpha_l) \approx \pi^2 \gamma^2 / J^2 \ln(\omega_l / J)\). Any time variation in \(\mu\) will lead to \(\delta > \delta_{min}\). In addition the deviation from adiabaticity \(\delta_{min}\) is quadratically increasing with the dephasing parameter \(\gamma_p\). Under different conditions when \(\mu\) is large this tendency can be reversed.

Another source of external noise is due to amplitude errors in the control of the frequency \(\omega(t)\). These errors are modeled by a Gaussian random process described by the Lindblad term: \(L_A(\mathbf{A}) \Rightarrow -\gamma_p \omega^2 [\mathbf{B}_1, [\mathbf{B}_1, \mathbf{A}]]\), where \(\gamma_p\) characterizes the amplitude noise. The equation of motion for the noise propagator \(U_3\) becomes:

\[
\frac{d}{dt} U_3(\Theta) = \mathcal{W}(\Theta) U_3(\Theta),
\]

(22)

We seek an approximate for small \(\mu\), using a similar procedure of calculating the propagator for one period \(U_3(\Theta = 2\pi)\). The \(U_3(1,1)\) element decouples from the remaining part of the propagator. As a result \(1 - \delta = U_3(1,1) \approx e^{-\gamma_p \omega^2 / \Omega^2 (\pi / \Omega)^2}(\Theta = 2\pi)\).

The smallest \(\delta\) is achieved for a one-period cycle, eq. (13), then \(\delta_{min} \approx \gamma_p J \frac{\omega^2}{\Omega^2} \).

Figure 2 shows \(\delta\) as a function of the propagation time for different values of \(\mu\) corresponding to the quantization condition, eq. (12), calculated numerically for phase (integrating eq. (18)) and amplitude noise (integrating eq. (22)). The oscillations reflect the periodicity of eq. (10), where frictionless solutions require an integer number of periods. The amplitude of the oscillations decreases with \(\mu\). The phase and amplitude noise have a different influence on the expansion dynamics.

Figure 3 shows the minimum temperature calculated numerically as a function of expansion time \(\gamma_{hc}\), for the phase and amplitude noise. The exact numerical calculation are consistent with the approximation when \(\mu \to 0\). The phase noise has a monotonic decrease of \(T_c(\min)\) reaching saturation as \(\gamma_{hc} \to \infty\), where \(T_c(\min) = -h \frac{1}{\sqrt{2 k_B \log(\gamma_p J / \gamma_{hc} / \sqrt{2})}}\). \(T_c(\min)\) of the amplitude noise is a monotonically increasing function of time which means that short expansion times lead to the minimum temperature.

If both amplitude and phase noise operate simultaneously the minimum temperature will be obtained at the crossing point. This optimum will move with the ratio \(\gamma_p / \gamma_a\).

Conclusions. — The necessary condition for the working medium to cool down to absolute zero is that \(\langle \mathbf{H} \rangle_0 = -h \mathbf{L}_0\). In order to start in the ground state at the hot end the working medium has to equilibrate with a very high frequency \(\Omega_0\). Perfect adiabatic following will maintain the system in its ground state which defines the frictionless solution. We found a family of additional frictionless solution obeying a quantization rule for \(\mu\), the adiabatic parameter.
The main result of this study is that any noise in the controls of \( \omega(t) \) will eliminate the frictionless solutions leading to a minimum temperature \( T_c(\text{min}) \). This finding is consistent with experiments on the demagnetization cooling of a gas [20] which obtained a minimum temperature an order of magnitude larger than the theoretical prediction [21], which attributes the discrepancy to the noise in the controls. The logarithmic dependence on the noise parameters means that \( T_c(\text{min}) \) is of the order of \( J \). The devastating effect of the amplitude noise is caused by the non-commutativity of \( \hat{H}_{\text{ext}}(\omega) \) with the instantaneous Hamiltonian \( \hat{H}(t) \). In this case the friction increases with the expansion time, i.e. the noise accumulates. The surprise is the negative effect of phase noise. Its generator, eq. (15), can be interpreted as the result of the weak measurement of \( \hat{H} \), forcing collapse to a state diagonal in energy [22]. It would seem that this should improve adiabaticity and reduce friction. For an engine model we found such an effect of enhanced performance [18]. At the minimum temperature the refrigerator dephasing is harmful. Dephasing destroys the periodic \( \delta = 0 \) solutions. On the contrary, the noise reduces \( \delta \) in the non-periodic expansions so it acts as a lubricant, nevertheless these expansions do not lead to minimum temperature. The optimal cooling conditions with minimum \( \delta \) are obtained for small and constant \( \mu \). Under these conditions \( \delta \) is dominated by the dephasing noise and \( \delta_{\text{min}} \) reaches a non-zero minimum asymptotic value when \( \tau_{hc} \to \infty \) and with \( \mu \to 0 \). This noise causes \( T_c(\text{min}) > 0 \), i.e. the absolute zero is not connected. In the thermodynamic tradition is it possible to generalize from this example? Preliminary analysis of other models which without noise show a frictionless expansion [7,23] support this ansatz.

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