

LETTERS

Thermodynamic control by frequent quantum measurements

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Heat flow between a large thermal ‘bath’ and a smaller system brings them progressively closer to thermal equilibrium while increasing their entropy¹. Fluctuations involving a small fraction of a statistical ensemble of systems interacting with the bath result in deviations from this trend. In this respect, quantum and classical thermodynamics are in agreement^{1–5}. Here we predict a different trend in a purely quantum mechanical setting: disturbances of thermal equilibrium between two-level systems (TLSs) and a bath⁶, caused by frequent, brief quantum non-demolition^{7–10} measurements of the TLS energy states. By making the measurements increasingly frequent, we encounter first the anti-Zeno regime and then the Zeno regime (namely where the TLSs’ relaxation respectively speeds up and slows down^{11–15}). The corresponding entropy and temperature of both the system and the bath are then found to either decrease or increase depending only on the rate of observation, contrary to the standard thermodynamical rules that hold for memory-less (Markov) baths^{2,5}. From a practical viewpoint, these anomalies may offer the possibility of very fast control of heat and entropy in quantum systems, allowing cooling and state purification over an interval much shorter than the time needed for thermal equilibration or for a feedback control loop.

To understand the origins of the predicted anomalies, consider a thermal bath in equilibrium with an ensemble of quantum systems. The energy of the quantum systems is briefly measured, with the following effects. Classically, the equilibrium state of the systems may remain intact, since measurements can be chosen to be non-intrusive, that is, to involve no energy exchange, and merely to provide ‘snapshots’ of the overall system. Likewise, quantum mechanically, nearly ideal (projective) measurements involve no energy cost when performed by macroscopic detectors on isolated systems⁸. However, finite-time coupling, followed by abrupt decoupling, of two quantum ensembles, which may be viewed as the detection of one ensemble by the other, may cause an increase in their mean total energy¹⁶. Controlled perturbations of open quantum systems may also cause their properties to change anomalously^{17,18}.

Here we are concerned with the time-resolved evolution of two-level systems (TLSs) that are initially in thermal equilibrium with a much larger bath. We investigate how the temperature and entropy of these systems evolve when the systems are probed by brief, repeated measurements. Our analysis shows that the temperature and entropy display a universal dependence on the time interval between measurements, provided that they are frequent enough to cause deviations from energy conservation in the system–bath exchange¹⁹, consistent with time–energy uncertainty.

Quantitatively, we consider a TLS with energy separation $\hbar\omega_a$ (‘a’ denoting ‘atom’) that is weakly coupled to a thermal bath of harmonic oscillators, characterized by a bath response or correlation (memory) time of $t_c \gg 1/\omega_a$, which typically marks the onset of equilibrium. After equilibrium has been reached, we perform $k = 1, \dots, K$ quantum

non-demolition measurements^{7–10} of the TLS energy states at times separated by $\Delta t_k = t_{k+1} - t_k$. Each measurement has a brief duration of $\tau_k \ll 1/\omega_a$. Our aim is to explore the evolution as a function of the time separation $\Delta t_k \leq 1/\omega_a \ll t_c$ between consecutive measurements in the uncharted non-Markov domain.

We show that this process entails three universal anomalies. The first of these is that the quantum mechanical non-commutativity of the system–detector and system–bath interactions causes the system to heat up immediately after the measurement, at the expense of the detector–system coupling but not at the expense of the bath. This heating up occurs only for very low values of Δt_k compatible with the quantum Zeno effect (QZE)¹¹. The second anomaly is that a transition of the TLS ensemble from heating to cooling may occur as we vary the interval between consecutive measurements from $\Delta t_k \ll 1/\omega_a$ to $\Delta t_k \approx 1/\omega_a \ll t_c$. This marks the transition from values of Δt_k compatible with the QZE to values compatible with the anti-Zeno effect (AZE)^{12–15}. We note that the cooling may occur even if the bath is initially hotter. The third anomaly is that, correspondingly, oscillations of the entropy relative to that of the equilibrium state take place, contrary to the markovian notion of the second law of thermodynamics^{2,5}.

This generic scenario is governed by the following total hamiltonian of the system that interacts with the bath and is intermittently perturbed by the coupling of the system to the detector (measuring apparatus):

$$H(t) = H_{\text{tot}} + H_{\text{SD}}(t)$$

$$H_{\text{tot}} = H_S + H_B + H_{\text{SB}}$$

Here H_S is the hamiltonian of the TLS, with ground and excited states $|g\rangle$ and $|e\rangle$, respectively; H_B is that of the thermal bath composed of harmonic oscillators with energies $\hbar\omega_\lambda$ (λ denoting the mode of the bath); $H_{\text{SB}} = \text{SB}$ is the system–bath interaction hamiltonian¹⁹ (the spin-boson interaction), which is a product of the system-dipole (or spin-flip) operator S and the operator B describing the bath excitations and de-excitations; and $H_{\text{SD}}(t)$ is the time-dependent measurement hamiltonian that couples the system to a detector consisting of energy-degenerate ancillae (for details, see Supplementary Information section A). It is essential that the coupling hamiltonians H_{SB} and H_{SD} do not invoke the rotating-wave approximation¹⁹; that is, we do not impose energy conservation between the system and the bath or the detector, on the timescales considered¹⁵.

The near-equilibrium state ρ_{tot} before a measurement has three pertinent characteristics (see Supplementary Information section B). First, it displays system–bath entanglement with off-diagonal matrix elements $\langle e|\rho_{\text{tot}}|g\rangle \neq 0$. Second, the system is described by a diagonal reduced density matrix $\rho_S = \text{Tr}_B \rho_{\text{tot}}$ (Tr_B denoting trace over the bath) in the H_S eigenbasis. Third, the mean interaction energy $\langle H_{\text{SB}} \rangle$ is negative, assuming that ρ_{tot} weakly deviates from the ground

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state of H_{tot} : $\langle H_{\text{SB}} \rangle = \langle H_{\text{tot}} \rangle - \langle H_{\text{S}} + H_{\text{B}} \rangle < 0$. This is because the correction to the ground-state energy of H_{tot} due to a weakly perturbing interaction H_{SB} is negative (to leading second order).

We next consider the disturbance of this equilibrium state by a nearly impulsive (projective) quantum measurement ($\tau \rightarrow 0$) of the TLS, in the basis $\{|g\rangle, |e\rangle\}$. Such measurements do not resolve the energies of the TLS states, due to time-energy uncertainty. However, they can discriminate between states of different symmetry, for example different angular momenta¹⁹. The measurement correlates the TLS energy eigenstates with mutually orthogonal states of an ancillary detector and has distinctly quantum mechanical consequences (see Supplementary Information section A): using the energy supplied by $H_{\text{SD}}(0 < t < \tau)$ (the system-detector coupling), but without changing $\langle H_{\text{D}} \rangle$, it eliminates the mean system-bath interaction energy; that is, it sets $\langle H_{\text{SB}} \rangle$, the pre-measurement value of which was negative, to zero. Thus, for $0 \leq t \leq \tau$:

$$\begin{aligned} \langle H_{\text{SB}}(0) \rangle < 0 &\rightarrow \langle H_{\text{SB}}(\tau) \rangle = 0 \\ \langle H_{\text{SD}}(t) \rangle &= -\langle H_{\text{SB}}(t) \rangle \end{aligned}$$

This energy transfer, resulting in a change in the entanglement between the system and the bath, triggers the quantum dynamics that redistributes their mean energy and entropy.

The information gained from such measurements may be used to sort system sub-ensembles according to their measured energy, to extract work or entropy change, in the spirit of Maxwell's demon²⁰. Here, however, we let the entire TLS ensemble evolve independently of the measured result; that is, we trace out the detector states, a procedure known as non-selective, or unread, measurements¹⁰.

After the measurement (because $H_{\text{SD}}(t \geq \tau) = 0$), time-energy uncertainty at $\Delta t \leq 1/\omega_a$ results in the breakdown of the rotating-wave approximation; that is, $\langle H_{\text{S}} + H_{\text{B}} \rangle$ is not conserved as Δt grows. Only $\langle H_{\text{tot}} \rangle$ is conserved, by unitarity, until the next measurement. Hence, the post-measurement decrease of $\langle H_{\text{SB}} \rangle$ with Δt , signifying the restoration of equilibrium, $\langle H_{\text{SB}}(\tau) \rangle = 0 \rightarrow \langle H_{\text{SB}}(\tau + \Delta t) \rangle < 0$, is at the expense of an increase in $\langle H_{\text{S}} + H_{\text{B}} \rangle = \langle H_{\text{tot}} \rangle - \langle H_{\text{SB}} \rangle$, that is, a heating of the system and the bath combined:

$$\begin{aligned} \frac{d}{dt} (\langle H_{\text{S}} \rangle + \langle H_{\text{B}} \rangle) \Big|_{\tau + \Delta t} &> 0 \\ \frac{d}{dt} \langle H_{\text{SB}} \rangle \Big|_{\tau + \Delta t} &< 0 \end{aligned} \quad (1)$$

The post-measurement evolution of the system alone, described by $\rho_{\text{S}} = \text{Tr}_{\text{B}} \rho_{\text{tot}}$, is not obvious. Its Taylor expansion:

$$\rho_{\text{S}}(\tau + \Delta t) \approx \rho_{\text{S}}(\tau) + \Delta t \dot{\rho}_{\text{S}}(\tau) + \frac{\Delta t^2}{2} \ddot{\rho}_{\text{S}}(\tau) + \dots \quad (2)$$

where dots denote differentiation with respect to time, holds at short evolution times $\Delta t \ll 1/\omega_a$. The zeroth-order term is unchanged by the measurement: $\rho_{\text{S}}(\tau) = \rho_{\text{S}}(t \leq 0)$. The first derivative vanishes at $t = \tau$ ($\Delta t = 0$) owing to the definite parity of the bath density operator correlated with $|g\rangle$ or $|e\rangle$ (see Supplementary Information section B). This initial post-measurement vanishing, $\dot{\rho}_{\text{S}}(\tau) = 0$, is the QZE condition^{11,13-15}. The time evolution of ρ_{S} is then governed by its second time derivative, $\ddot{\rho}_{\text{S}}(\tau)$, which can be shown (see Supplementary Information section B) to have the same sign as $\sigma_z = |e\rangle\langle e| - |g\rangle\langle g|$, the z Pauli matrix, which is the population difference operator of the TLS. Hence, the second derivative in equation (2) is positive shortly after the measurement, consistent with equations (1), if there is no initial population inversion of the system, that is, for non-negative temperature.

The evolution of ρ_{S} at longer times (in the regime of weak system-bath coupling) may be approximately described (as verified by our exact numerical simulations^{21,22}; see Supplementary Information section C) by the second-order non-markovian master equation (Fig. 1 main panel). Because ρ_{S} is diagonal, the master equation describing it can be cast into the following population-rate equations¹⁵ (we drop the subscript 'S' in what follows and set the measurement time to be $t = 0$):

$$\dot{\rho}_{ee}(t) = -\dot{\rho}_{gg}(t) = R_g(t)\rho_{gg} - R_e(t)\rho_{ee} \quad (3)$$

$$R_e(t) = 2t \int_{-\infty}^{\infty} d\omega G_{\text{T}}(\omega) \text{sinc}[(\omega - \omega_a)t] \quad (4)$$

$$R_g(t) = 2t \int_{-\infty}^{\infty} d\omega G_{\text{T}}(\omega) \text{sinc}[(\omega + \omega_a)t]$$

$$G_{\text{T}}(\omega) = (n_{\text{T}}(\omega) + 1)G_0(\omega) + n_{\text{T}}(-\omega)G_0(-\omega)$$

Here $\text{sinc}(x) = \sin(x)/x$ is the Fourier transform of the time interval after the measurement, $G_{\text{T}}(\omega)$ is the temperature-dependent coupling spectrum of the bath, $G_0(\omega)$ is the zero-temperature coupling spectrum with peak coupling strength γ at ω_0 and spectral width $\sim 1/t_c$, and $n_{\text{T}}(\omega) = (e^{\beta\omega} - 1)^{-1}$ is the inverse-temperature-dependent (that is, β -dependent) population of bath mode ω .

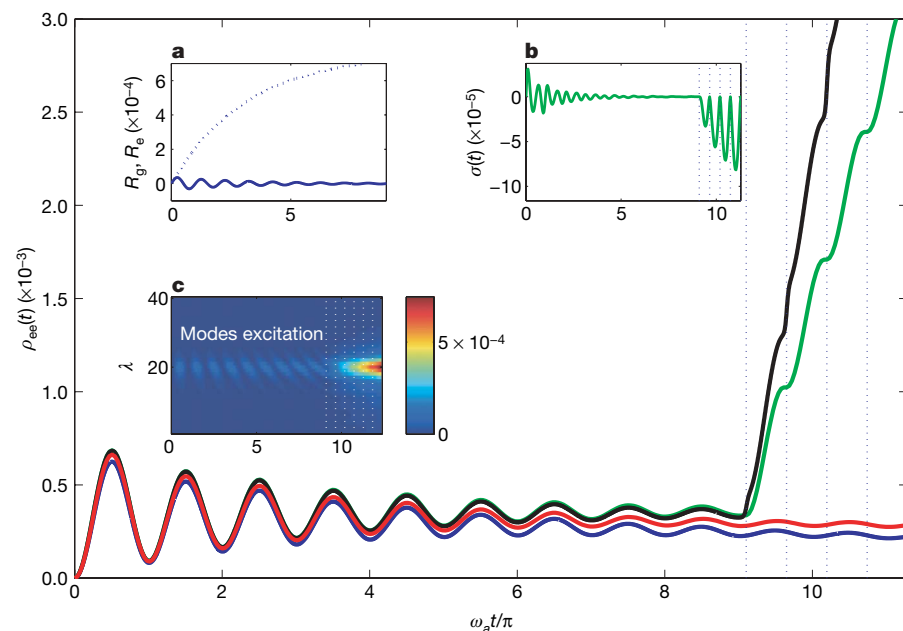


Figure 1 | System and bath evolution as a function of time. The main panel shows the excited-level population as a function of time for the initial zero-temperature product state of the system and bath. The population relaxes to quasi-equilibrium after a few oscillations. It is then subjected to a series of measurements (vertical dotted lines). Measurements of finite duration ($\tau_k = 0.11/\omega_a$) (black line) result in a larger heating up than do impulsive measurements (green line), but the dominant effect is the same for both. We observe agreement between the results of the second-order master equation (blue line), two-quanta exchange with a discrete bath (black and green lines), and the exact numerical solution for a discrete bath of 40 modes (red line). The horizontal axes of the inset panels all indicate $\omega_a t / \pi$, as in the main panel: **a**, relaxation rates R_g (solid blue line) and R_e (dotted line) as functions of time; **b**, negative of the rate of change of relative entropy, $\sigma(t)$; **c**, excitations of the 40 bath modes in the two-quanta model as functions of time. Parameters: $t_c = 10/\omega_a$, $\omega_0 = \omega_a$, $\gamma = 0.07\omega_a$.

The dynamics of equation (3) is determined by $R_e(t)$ and $R_g(t)$, which are respectively the relaxation rates of the excited and ground states. Their non-Markov time dependence yields three distinct regimes, universally dependent on the post-measurement times t of the spin-boson evolution.

First, at short times $t \ll 1/\omega_a \ll t_c$, the sinc functions in equations (4) are spectrally much broader than is G_T . The relaxation rates R_e and R_g are thus equal at any temperature, indicating the complete breakdown of the rotating-wave approximation discussed above: the $|g\rangle \rightarrow |e\rangle$ and $|e\rangle \rightarrow |g\rangle$ transitions do not require quantum absorption or emission by the bath, respectively. The rates R_e and R_g thus become linear in time, manifesting the QZE^{13–15}:

$$R_e(t \ll t_c) \approx 2\dot{R}_0 t$$

$$R_g(t \ll t_c) \approx 2\dot{R}_0 t$$

$$\dot{R}_0 \equiv \int_{-\infty}^{\infty} d\omega G_T(\omega) = \langle B^2 \rangle$$

This implies that in the short-time regime we have the universal Zeno heating rate:

$$\frac{d}{dt} (\rho_{ee} - \rho_{gg}) \approx 4\dot{R}_0 t (\rho_{gg} - \rho_{ee})$$

Second, at intermediate non-markovian times $t \approx 1/\omega_a$, when the sinc functions and G_T in equations (4) have comparable widths, the relaxation rates R_e and R_g exhibit several unusual phenomena that stem from time–energy uncertainty. The change in the overlap of the sinc and G_T functions with time results in damped aperiodic oscillations of $R_e(t)$ and $R_g(t)$, near the frequencies $\omega_0 - \omega_a$ and $\omega_0 + \omega_a$, respectively. We call this oscillatory time dependence, which is displayed in neither the QZE nor the converse AZE of accelerated relaxation^{12–14}, the oscillatory Zeno effect (OZE). Owing to the negativity of the sinc function between its consecutive maxima, we can have a negative relaxation rate,

which is completely forbidden by the rotating-wave approximation. Because $\text{sinc}[(\omega + \omega_a)t]$ is shifted much further from the peak of $G_T(\omega)$ than is $\text{sinc}[(\omega - \omega_a)t]$, $R_g(t)$ is more likely to be negative than is $R_e(t)$ (Figs 1a, 2a). Hence, $\rho_{gg}(t)$ may grow at the expense of $\rho_{ee}(t)$ more than is allowed by the thermal–equilibrium detailed balance. This may cause transient cooling, as described below.

Third, at long times $t \gg t_c$, the relaxation rates attain their ‘golden rule’ (Markov) values¹⁵:

$$R_e(t \gg t_c) \approx 2\pi G_T(+\omega_a)$$

$$R_g(t \gg t_c) \approx 2\pi G_T(-\omega_a)$$

The populations thus approach those of an equilibrium Gibbs state the temperature of which is equal to that of the thermal bath (Fig. 1 main panel).

We now consider entropy dynamics. The entropy of ρ_S relative to its equilibrium state ρ_0 (‘entropy distance’), and the negative of its rate of change, can always be defined^{3,5}:

$$S(\rho_S(t) || \rho_0) \equiv \text{Tr}\{\rho_S(t) \ln \rho_S(t)\} - \text{Tr}\{\rho_S(t) \ln \rho_0\} \quad (5)$$

$$\sigma(t) \equiv -\frac{d}{dt} S(\rho_S(t) || \rho_0)$$

In the markovian realm, σ thus defined is identified as the ‘entropy production rate’^{2,3,5}, and $\sigma(t) \geq 0$ is a statement of the second law of thermodynamics in this realm. Because ρ_S is diagonal, it follows (see Supplementary Information section D) that $\sigma(t)$ is positive if and only if $d(|\rho_{ee}(t) - (\rho_0)_{ee}|)/dt \leq 0$, consistent with the interpretation of the relative entropy $S(\rho_S || \rho_0)$ in equations (5) as the entropic distance from equilibrium. However, in the short-time regime, whenever the oscillatory $\rho_{ee}(t)$ drifts away from its initial or final equilibrium σ takes negative values (Fig. 1b) that are forbidden by the markovian statement of the second law.

To realistically model the repeated measurements, that is, give them finite duration, we assume a smooth temporal profile of the

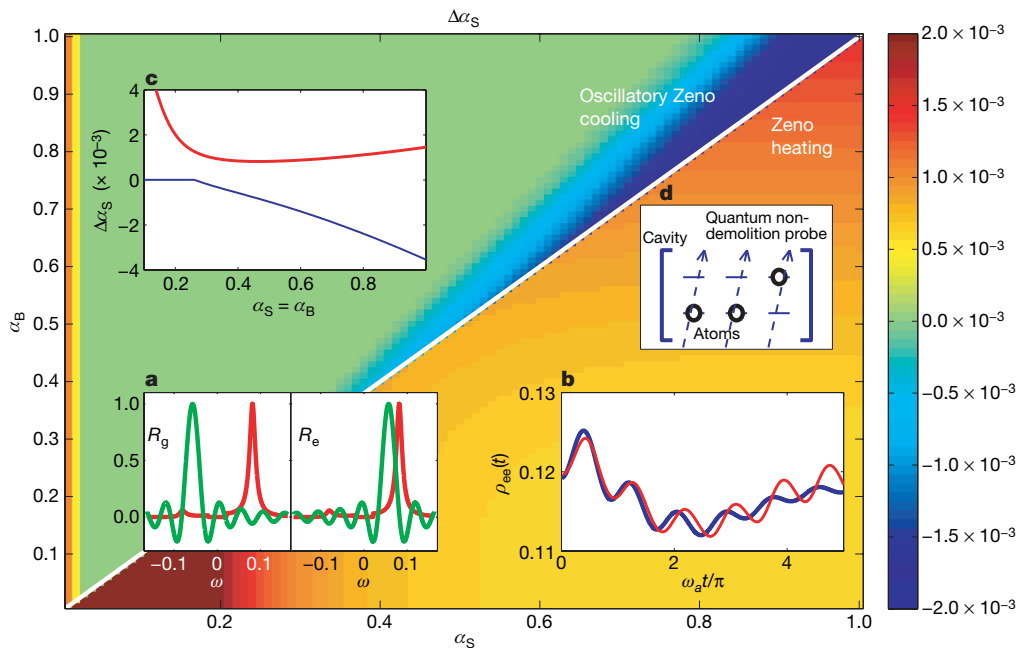


Figure 2 | Maximal heating and cooling of the system. The main panel shows the maximal heating (lower half) and cooling (upper half) of the system, for different system (horizontal axis) and bath (vertical axis) initial temperatures α_S and α_B ($\alpha_S \neq \alpha_B$), where $\alpha_S = \hbar\omega_a/\beta_S$, $\alpha_B = \hbar\omega_a/\beta_B$, and $\Delta\alpha_S = \max(\alpha_S(t) - \alpha_S(0))$ for heating and $\Delta\alpha_S = \min(\alpha_S(t) - \alpha_S(0))$ for cooling. (β denotes the inverse temperature.) **a**, R_g and R_e (see equations (4)), depicted as spectral overlaps of $G_T(\omega)$ (red) and $\text{sinc}[(\omega + \omega_a)t]$ (for R_g) and $\text{sinc}[(\omega - \omega_a)t]$ (for R_e) (green). **b**, Example of a system undergoing

first Zeno heating and then oscillatory Zeno cooling, obtained from the second-order master equation (blue) and from the exact numerical solution for a discrete bath of 40 modes (red). **c**, Maximal Zeno heating (red) and subsequent maximal oscillatory Zeno cooling (blue), as functions of the common initial temperature of the system and the bath $\alpha_S = \alpha_B$. We note the presence of a critical temperature for oscillatory Zeno cooling. Parameters: $t_c = 10/\omega_a$, $\omega_0 = \omega_a/0.7$, $\gamma = 4.36\omega_a$. These effects can be strongly magnified by choosing other suitable parameters. **d**, Possible experimental setup.

coupling to the detector (see Supplementary Information section A). The k th measurement then occurs at time t_k and has a duration of τ_k . Figure 1 (main panel) compares the population evolution using projective (impulsive) and finite-duration measurements with $\tau_k \approx 0.1/\omega_a$. Finite-duration measurements increase the Zeno heating in comparison with impulsive ones on account of the extra energy supplied by the coupling to the detector. However, the basic effect is seen (Fig. 1 main panel) to be the same and is governed by the time derivative of $\langle H_{SB} \rangle$ (see equations (1)). Counterintuitively, finite-duration measurements can increase the cooling, despite the extra energy supplied by the apparatus.

If we repeat this procedure often enough, the TLSs will increasingly either heat up or cool down, if we choose the time intervals Δt_k to coincide with either peaks or troughs of the ρ_{ee} oscillations, respectively. The minimal value of σ can also be progressively lowered with each measurement (Fig. 1b). Because consecutive measurements affect the bath and the system differently, they may acquire different excitations or de-excitations, which then become the initial conditions for subsequent QZE heating or OZE cooling. The results are shown in Fig. 2 for both different (main panel) and common (Fig. 2c) initial temperatures of the system and the bath. We note that the system may heat up on account solely of the QZE, although the bath is initially colder, or cool down on account solely of the OZE or AZE, although the bath is initially hotter (Fig. 2 main panel). The bath may also undergo changes in energy and entropy (Fig. 1c).

One experimental realization of these effects can involve atoms or molecules in a microwave cavity (Fig. 2d) with a coupling spectrum $G_T(\omega)$ (controllable using the cavity quality factor and temperature) centred at ω_0 . Measurements can be effected on such a TLS ensemble with resonance frequency ω_a in the microwave domain, at time intervals $\Delta t_k \approx 1/(\omega_0 \pm \omega_a)$, using an optical quantum non-demolition probe⁸ at a frequency ω_p such that $\omega_p \gg \omega_a$ and $\omega_p \gg \omega_0$. The probe pulses undergo a different Kerr-nonlinear phase shift $\Delta\phi_e$ or $\Delta\phi_g$ depending on the different symmetries (for example angular momenta) of $|e\rangle$ and $|g\rangle$. The relative abundance of $\Delta\phi_e$ and $\Delta\phi_g$ thus reflects the ratio $\rho_{ee}(t_k)/\rho_{gg}(t_k)$. Such quantum non-demolition probing may be performed with time duration much shorter than ω_a^{-1} , that is, $\omega_a\tau_k \ll 1$, without resolving the energies of $|e\rangle$ and $|g\rangle$. Ion traps or solid matrices may serve as phonon baths instead of photonic cavities.

Since non-selective measurements increase the von Neumann entropy of the detector ancillae, their entropic price precludes a 'perpetuum mobile' if closed-cycle operation is attempted. However, if our ancillae are laser pulses, they are only used once and we may progressively change the TLS ensemble thermodynamics by consecutive pulses, disregarding their entropic or energetic price. The practical advantage of the predicted anomalies is the possibility of very rapid control of cooling and entropy, which may be attained after several measurements at $t \geq \omega_a^{-1}$ and is only limited by the measurement rate. By contrast, conventional cooling requires much longer times, $t \gg t_c$, to reach thermal equilibrium. Likewise, temperature control based on a feedback loop is inherently more time consuming. The proposed fast cooling should be advantageous for quantum information processing and storage based on quantum state distillation.

The findings that we report here establish a new link between frequent quantum measurements, which are the operational probes of short-time evolution, and non-equilibrium thermodynamical anomalies: heat and entropy rates of change that have the 'wrong' sign, rather than displaying their usual monotonic approach to equilibrium. These anomalies are determined by the oscillatory or negative values of the non-markovian quantum relaxation rates at short times corresponding to large energy uncertainty. They reveal unfamiliar general aspects of post-measurement quantum dynamics: the AZE, which was initially proposed as a means of enhancing or accelerating the initial-state change^{12–15}, here can either restore the equilibrium

state or further depart from it via cooling. These anomalies underscore the fact that the system and the bath are inseparable (entangled)^{6,23}, even under weak-coupling conditions, a fact that has important implications for their short-time dynamics.

These results prompt further studies of a hitherto unexplored non-markovian short-time domain where neither the existing formulations of the second law of thermodynamics^{2,5} nor the common notion that heat always flows from hotter to colder ensembles is applicable. This calls for an examination of non-markovian designs of quantum heat engines and their comparison to their markovian counterparts²⁴. The short-time domain may also necessitate an in-depth scrutiny of fundamental quantum thermodynamical concepts. In particular, the need for temporal 'coarse graining' of entropy should be examined.

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Supplementary Information is linked to the online version of the paper at www.nature.com/nature.

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