

Erasure of a nonequilibrium memory beyond Landauer's bound using levitated optomechanics with spatio-temporal optical control

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Landauer's principle provides a deep connection between information processing and thermodynamics by setting a lower limit on the energy consumption and heat production of logically irreversible transformations. However, Landauer's original formulation assumes that information is initially stored in an equilibrium state, whereas real devices often operate with out-of-equilibrium states. Here, we experimentally demonstrate that nonequilibrium memory states can provide a useful resource for fast information erasure below the equilibrium bounds by using an underdamped levitated optomechanical system. To that end, we introduce a versatile optical levitation scheme that enables fast and precise dynamical shaping of general potential landscapes. By harnessing the energy and entropy of an initial nonequilibrium two-state memory, we demonstrate reduced power consumption as well as negative heat production during erasure. Our findings significantly broaden the range of tools available in levitodynamics and suggest that the engineering of appropriate far-from-equilibrium memory states could pave the way for an approach to heat management.

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I. INTRODUCTION

Heat production poses crucial challenges for computer hardware development. From a physical perspective, dissipation in logically irreversible operations is tightly connected to Landauer's bound [1], a central result of information thermodynamics [2–5]. Specifically, Landauer's principle, in its original equilibrium formulation, states that the erasure of a single bit of information in an environment at temperature T (Fig. 1) consumes at least $kT \ln 2$ of work and entails the same amount of dissipated heat (k denotes the Boltzmann constant) [1–5]. Today, real computational devices operate several orders of magnitude above this value [6], but they are expected to approach it in the next few decades [7]. Since the second law of equilibrium thermodynamics sets fundamental lower limits to heat dissipation during erasure, nonequilibrium strategies need to be explored to improve heat management in these devices.

Meanwhile, proof-of-principle experiments with colloidal particles, nanomagnetic systems, and micromechanical oscillators have already achieved irreversible information erasure with an energy dissipation close to the Landauer limit [8–12]. In particular, Landauer's bound has been reached for

near-equilibrium erasure times on the order of minutes for overdamped colloidal particles trapped in a double-well potential [8,10] and, lately, in only 100 ms with an underdamped cantilever [12]. Beyond such confirmations, the need to extend Landauer's principle for deviations from an idealistic near-equilibrium scenario has been recognized: Implementations with finite success probability [13,14], asymmetric memories [15–17], and finite time [18–21] have, for instance, been studied both theoretically and experimentally. Recent theoretical work has further shown the significant impact of the engineering of far-from-equilibrium states on the thermodynamic bounds of erasure [22–24], and emphasized the necessity to generalize Landauer's principle. In addition, real-world memory devices, such as modern volatile memories [25], store information in nonequilibrium states that cost energy and entropy to create, and are lost when power is turned off. Despite their widespread use, the impact of far-from-equilibrium *initial memory states* on Landauer's bound has yet to be explored experimentally.

Here, we use an underdamped levitated optomechanical system to store information in a nonequilibrium state of a tunable double-well potential. Such initial memory states are predicted to allow control of both work consumption and heat dissipation, which can be reduced significantly below the equilibrium limits by harnessing energy and entropy of the prepared state [22–24]. Exploiting this property of the initial state contrasts with the investigation of nonequilibrium erasure processes, which typically cause additional dissipation compared to equilibrium protocols (see, e.g., [18–21]). It is important to emphasize the distinction between both

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sources of nonequilibrium effects (nonequilibrium state versus nonequilibrium erasure process), as the erasure process in our experiment is, by design, mostly in equilibrium. The improvement over Landauer's bound here is based on the nonequilibrium initial memory state (NIMS), whose nonzero entropic distance to equilibrium constitutes a useful resource for information erasure. Optical levitation offers a unique experimental platform for preparing nonequilibrium initial states and implementing thermodynamic protocols [26]. It enables control over the coupling between the trapped mechanical system and its environment, while also allowing flexible tuning of the optical potential [27]. These capabilities have been exploited in various experiments, including fast modulation of Gaussian optical traps, electrostatic trapping, and the static implementation of nonlinear potentials [28–34]. However, levitated mechanical systems do not currently offer the same level of control known from optical micromanipulation of overdamped colloids in liquid [35,36]. In particular, the implementation of the Landauer protocol in optical levitation requires even better controllability than available today, such as fast spatio-temporal manipulation, i.e., the shaping of nonlinear optical potentials with fast dynamical modulation, which has been lacking until now from the toolbox of levitodynamics [26,32,37–39]. The emulation of virtual potentials with feedback control, which has been employed in underdamped clamped systems [40], could in principle also be applied to optical levitation setups.

We go beyond these results by demonstrating dynamically programed control of *real* anharmonic potentials via the spatial intensity distribution and an electrostatic tilting force to modulate generic nonlinear potentials on fast timescales in the underdamped regime. This versatile approach may also provide access to non-Gaussian states in the quantum regime [41–44]. In the context of the present manuscript, this technique has the advantage of enabling the implementation of the Landauer protocol without knowledge of the spatial position of the mechanical system representing the memory.

We use our experimental platform to prepare nonequilibrium memory states and demonstrate a reduction of both consumed work and dissipated heat during erasure far below the equilibrium value, even reaching negative values, in agreement with a generalized nonequilibrium Landauer bound [24]. Specifically, we consider thermodynamic cycles consisting of a preparation, an erasure, and a reinitialization phase. We first prepare a nonequilibrium initial memory state (characterized by a nonequilibrium parameter ϵ) to store information in a double-well potential using a levitated silica nanosphere confined in an optical trap [Fig. 2(a)]. This generic model for a one-bit memory encodes state “0” (“1”) when the particle is in the left (right) potential well. Information is erased by setting the memory to state “0” irrespective of the initial state [1–5]. We implement the erasure by modulating the shape of the potential by varying the power of two cross-polarized laser beams, decreasing the barrier height, and applying a tilt to the left [Fig. 2(a)]. We design the erasure protocol to minimize applied work and dissipated heat, using temporal and spatial control over the system parameters [Figs. 2(b) and 2(c)]. We specifically exploit optical levitation to control the strength of the underdamped coupling of the nanoparticle to its environment, achieving a reduction of the erasure time by

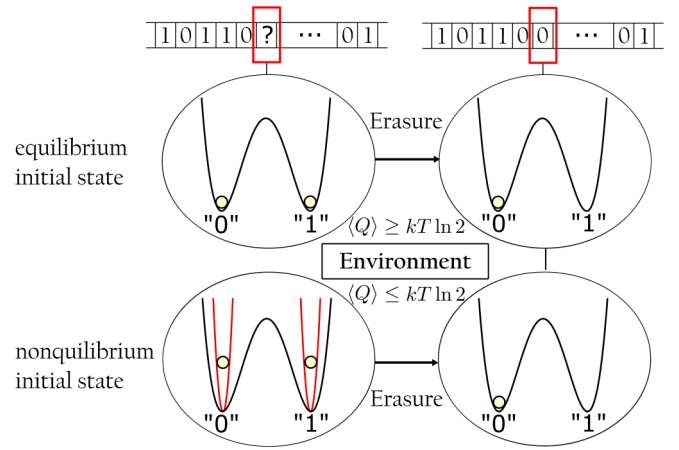


FIG. 1. Erasure of a nonequilibrium memory state. One bit of information stored in a generic, symmetric double-well memory may be erased by setting the system to state “0,” irrespective of the initial state “0” or “1.” While heat dissipated into a bath at temperature T is always larger than $kT \ln 2$ when information is stored in an equilibrium state, it might be smaller when energy and entropy of the initial nonequilibrium memory state are properly harnessed.

5 orders of magnitude compared to overdamped experiments [8,10,17].

II. GENERALIZATION OF LANDAUER'S BOUND FOR NONEQUILIBRIUM INITIAL MEMORY STATES

The implementation of our experimental protocol is divided into three steps: (1) a “preparation phase” during which we first prepare the nonequilibrium initial memory state, (2) an “erasure phase” during which we perform the erasure protocol, and (3) a “reinitialization phase” during which we reinitialize the particle to be equilibrated to the initial symmetric potential, to complete the cycle [Figs. 2(b) and 2(c)]. The preparation phase is absent for equilibrium memories as considered initially by Landauer [1–5]. Because of the initial preparation of a nonequilibrium memory state at the beginning of an erasure cycle, Landauer's principle needs to be extended [22]. Applying the laws of thermodynamics to this nonequilibrium configuration, a generalized Landauer bound holds for the mean dissipated heat $\langle Q \rangle$ and the mean consumed work $\langle W \rangle$ during the erasure phase [24],

$$\langle Q \rangle \geq Q_L = T \Delta I - \Delta \langle U_{\text{res}} \rangle, \quad (1)$$

$$\langle W \rangle \geq W_L = T \Delta I, \quad (2)$$

where $\Delta \langle U_{\text{res}} \rangle = \langle W \rangle - \langle Q \rangle$ is the variation of internal energy during an erasure (information reset) of duration t_{er} and $\Delta I = I(t_{\text{eq}}) - I(0)$ is the change of relative entropy, $I(t) = k \iint \rho(x, v, t) \ln[\rho(x, v, t)/(\rho_{\text{eq}}(x, v, t))] dx dv$, an entropic distance between a state $\rho(x, v, t)$ and the corresponding equilibrium state $\rho_{\text{eq}}(x, v, t)$ at time t (x is the position and v the velocity of the particle) [45] (Appendix A). The brackets $\langle \cdot \rangle$ denote an ensemble average over many process repetitions. Equations (1) and (2) reduce to the standard Landauer limits for symmetric memories and initial/final equilibrium states that correspond to $\Delta I = I(t_{\text{er}}) = k \ln 2$ and $\Delta \langle U_{\text{res}} \rangle = 0$. The

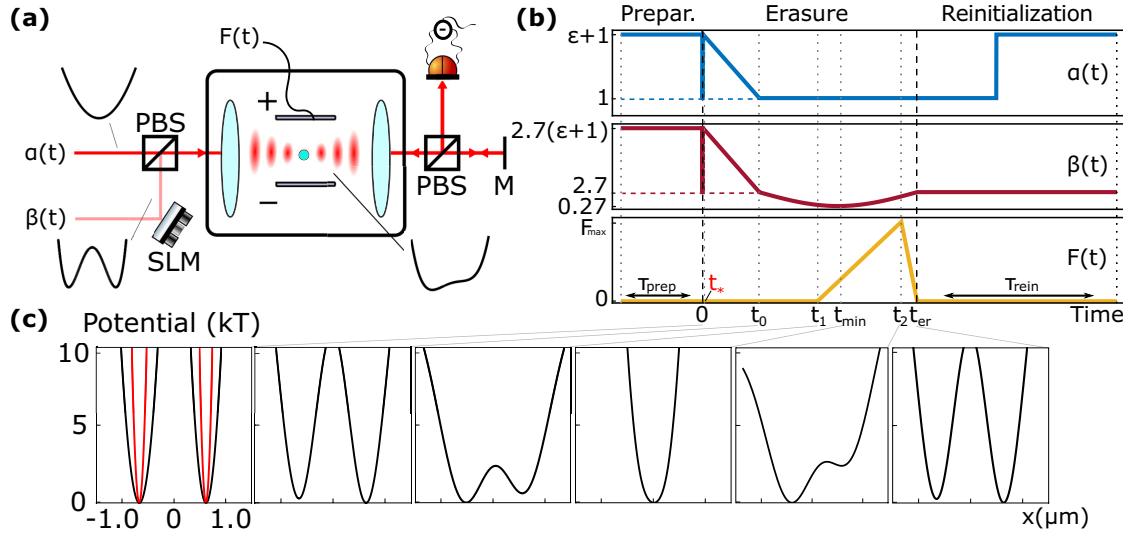


FIG. 2. (a) Optical levitation with spatio-temporal control. Two laser beams, one in a horizontally polarized TEM_{00} mode and the other prepared in a vertically polarized TEM_{10} mode using a spatial light modulator (SLM), are combined at a polarizing beam splitter (PBS) and tightly focused inside a vacuum chamber to create an optical trap. The power of the two laser beams, $\alpha(t)$ and $\beta(t)$, can be varied independently by acousto-optical modulators to modulate the shape of the double-well potential. Electrodes are placed near the optical trap (distance $100\ \mu\text{m}$) to control the tilt $F(t)$. After the vacuum chamber, the major part of the TEM_{00} mode is reflected back into the tweezer using the mirror M, creating a standing wave trap. A small fraction of the horizontally polarized light from the optical trap is fed into a standard split mirror detection to measure the position of the charged nanoparticle in the radial direction. (b) Experimental erasure protocol. Time dependence of the three control parameters $\alpha(t)$, $\beta(t)$, and $F(t)$ during the protocol (solid). The first section (preparation phase, duration $\tau_{\text{prep}} = 500\ \mu\text{s}$) prepares the nonequilibrium initial memory state by first letting the particle thermalize in a high-power potential and then switching back to the low-power memory potential. At $t = 0$, the erasure protocol starts with that setting. After erasure, a reinitialization phase ensures that the initial memory state is initialized repeatedly to the same low-intensity potential (duration $\tau_{\text{rein}} = 1250\ \mu\text{s}$); the yaxes are not to scale. The dashed lines show the equilibrium protocol (corresponding to $\epsilon = 0$). (c) Potential snapshots during erasure. Reconstructed trapping potentials at times, 0 (black) and $t_* \leq 1$ (red), $t_0 = 358$, $t_1 = 728$, $t_{\text{min}} = 858$, $t_2 = 1258$, and $t_{\text{er}} = 1358\ \mu\text{s}$. The red curve shows the potential for the nonequilibrium parameter $\epsilon = 4$.

generalized Landauer bound for heat Q_L may thus be tuned via the initial entropic distance to equilibrium $I(0)$ and the nonequilibrium preparation energy, $\Delta\langle U_{\text{pre}} \rangle = -\Delta\langle U_{\text{res}} \rangle$. On the other hand, the generalized Landauer bound for work W_L can be controlled through the entropic distance of the initial state to equilibrium $I(0)$. We note that the Landauer limit is less than $kT \ln 2$ for asymmetric memory states that store less than one bit of information [1–5].

III. OPTICAL POTENTIAL CONTROL: EXPERIMENTAL SETUP

In our experiment, we realize a dynamical double-well optical trap with the configuration shown in Fig. 2(a). A charged silica particle (radius $r = 71.5\ \text{nm}$) is confined in an optical tweezer inside a vacuum chamber. The pressure $P = (30.0 \pm 0.3)\ \text{mbar}$ is set to create $\Gamma = 2\pi \times (32.0 \pm 0.2)\ \text{kHz}$ of viscous damping. This value is close to the regime where the characteristic Kramer's time of the double well is significantly reduced [32]. In contrast to overdamped experiments [8,10,17], we thus operate at significantly reduced damping rates, and thus much shorter erasure times ($t_{\text{er}} \sim 1000\ \mu\text{s}$). The optical trap consists of two tightly focused (waist $W \approx 0.6\ \mu\text{m}$ and wavelength $\lambda = 1064\ \text{nm}$) and orthogonally polarized lasers with spatial modes TEM_{00} and TEM_{10} . The TEM_{00} mode creates a standard harmonic trap, with typical radial frequencies of around $2\pi \times 150\ \text{kHz}$ and an underdamped Q

factor. The TEM_{10} mode is obtained by using an SLM. We control their respective powers, $\alpha(t)$ and $\beta(t)$, independently using acousto-optic modulators (AOMs). Depending on their power ratio, a harmonic, quartic, or double-well trap with tunable barrier height is created transversely to the optical path along the x axis [Fig. 2(c)]. Our setup allows us to treat the motion along this axis independently of the other directions (Appendixes B and C). We further tilt the potential for the charged nanoparticle with a well-controlled force $F(t)$ generated by two electrode blades separated by $100\ \mu\text{m}$ along the x axis, mounted far outside the optical trap (Appendix B). The resulting potential is of the form $V(x, t) = -[\alpha(t)a + \beta(t)b/2x^2] \exp[-(cx^2)/2] + F(t)x$, where a , b , and c are geometric parameters that depend on the setup (Appendix D). This double-well potential models the two-state memory. The snapshots of the potential shown in Fig. 2(c) are generated by interpolating fits of independent measurements of the potential at equilibrium at various laser powers and electrode voltages and then interpolated to have a smooth function depending on the protocol time (Appendix D).

We create the NIMS in the double-well potential by first letting the particle thermalize in the potential with enhanced power $(\epsilon + 1)\alpha$ and $(\epsilon + 1)\beta$, during time τ_{prep} , before rapidly decreasing the power to $(\epsilon + 1)\alpha \rightarrow \alpha$ and $(\epsilon + 1)\beta \rightarrow \beta$ at $t = 0$ [Fig. 2(b), red line]. We label these states by a parameter ϵ that controls their departure from equilibrium. The rapid potential switch at the beginning of

the protocol takes approximately $1\ \mu\text{s}$ and is thus quasi-instantaneous. This timescale is faster than the damping rate and the system dynamics, ensuring that the particle momentum remains unchanged during the switch (see Appendix E). These parameters are chosen to decrease the energy and entropy of the initial state as compared to the equilibrium state given by $\epsilon = 0$. In particular, the nonequilibrium initial memory state has a narrower spatial distribution in both wells than the corresponding equilibrium state [i.e., it is thermally squashed in both wells (see, e.g., [46])]. Consequently, the distance to equilibrium increases with ϵ for any given memory potential shape. For a more detailed quantitative discussion of the states' deviation from equilibrium via the initial relative entropy $I(0)$ (see Appendix F). Note that other approaches, such as parametric modulation of the laser intensity or feedback control, allow preparation of the initial memory states prepared here [26,46,47].

For initial equilibrium states ($\epsilon = 0$), we maintain a power of 50 mW in the vacuum chamber in the TEM₀₀ mode at the start of the protocol and approximately 135 mW in the TEM₁₀ mode. This is the configuration with the lowest power to get a stable particle trap and a barrier high enough to avoid spontaneous jumping between the two wells. The nonequilibrium initial memory state is created by maintaining the same power proportion between the two beams and increasing the power according to ϵ . For example, the value $\epsilon = 4$ requires a power of 250 mW in the TEM₀₀ mode and of 675 mW in the TEM₁₀ mode. The reinitialization phase at the end of the protocol is moreover based on the random drop of the nanoparticle in either of the two wells before the protocol returns to the preparation phase.

We resolve the dynamics $x(t)$ of the nanoparticle in the double-well potential using a split detection of the TEM₀₀ mode in transmission [28]. The geometric configuration of the trap and the polarization of the light imply that the two radial frequencies are in the range of 120 kHz (with the x frequency consistently lower than the y frequency) and the axial frequency z on the order of 50 kHz. Unlike experiments where the particle's motion is strictly harmonic, and thus the detector's response can be filtered around the relevant frequency, here the potential's time variation and anharmonicity prevent us from high-pass filtering the data. For this reason, we create a standing-wave trap along the axial (z) direction by back-reflecting part of the TEM₀₀ mode into the trap, effectively increasing the z frequency up to around 800 kHz [Fig. 2(a)]. We then perform the thermodynamic analysis of the data by filtering the time trace with a low-pass filter at 280 kHz, removing the residuals of the harmonic y and z motion. Detector and force calibration account for position-dependent sensitivity and are discussed in detail in Appendices B and C.

Given the potential model $V(x, t)$, we furthermore evaluate the stochastic work, $W = \int_0^{t_{\text{er}}} dt \partial V(x, t) / \partial t \rightarrow \sum_{i=0}^{t_{\text{er}}} [V(x_i, t_{i+1}) - V(x_i, t_i)]$, along individual trajectories $x(t)$ [48] by summing over the discrete changes in the potential energy for incremental time steps of $dt = 0.16\ \mu\text{s}$. We additionally determine the stochastic heat via the first law as $\Delta U = U(t_{\text{er}}) - U(0) = W - Q = m[v(t_{\text{er}})^2 - v(0)^2]/2 + [V(x(t_{\text{er}}), t_{\text{er}}) - V(x(0), 0)]$. We calculate the respective average values of work $\langle W \rangle$ and heat $\langle Q \rangle$ by repeating each protocol about 20 000 times.

IV. ERASURE PROTOCOL: RESULTS

To put our results into perspective, we first implement the erasure protocol for a memory initialized in an equilibrium state, $\rho(x, 0) = \exp(-V(x)/kT)/Z$ [8,10], where Z is the partition function. This instance corresponds to an initial parameter $\epsilon = 0$. We reconstruct the potential from the measured trajectories (Appendix D). We find the probability $p_{\epsilon=0}^0(0) = 36.6\%$ ($p_{\epsilon=0}^1(0) = 63.4\%$) for the nanoparticle to be in state “0” (“1”), indicating that the memory state is slightly asymmetric, and thus stores less than one bit of information. We take this asymmetry explicitly into account in the thermodynamic analysis (Appendix A). Using the equilibrium protocol, we set the memory to state “0” with a success probability $p_{\epsilon=0}^0(t_{\text{er}}) = 99.6\%$. The particle hence ends up in the unwanted state “1” in only 0.4% of all runs. The erasure is achieved by decreasing the barrier height via $\beta(t)$, while keeping the barrier width controlled by $\alpha(t)$ constant, and applying a tilt $F(t)$ (red and blue dashed lines in Fig. 2(b)). We obtain the mean applied work $\langle W \rangle_{\text{eq}} = (0.583 \pm 0.046)\text{ kT}$ and the mean dissipated heat $\langle Q \rangle_{\text{eq}} = (0.620 \pm 0.064)\text{ kT}$. These two values are close to the theoretical equilibrium Landauer bounds for work and heat, $\langle W \rangle_L^{\text{eq}} = \langle Q \rangle_L^{\text{eq}} = (0.7 \pm 0.10)\text{ kT}$, determined for the measured values of the probabilities $p_{\epsilon=0}^0(0)$ and $p_{\epsilon=0}^0(t_{\text{er}})$, for the asymmetric memory. For a nonequilibrium memory state, the above equilibrium erasure protocol (corresponding here to the protocol for an initial value $\epsilon \neq 0$) does not harness the preparation energy and entropy of the initial state: both are dissipated, leaving the two Landauer limits unchanged [24]. Figure 2(b) (solid blue, red, and yellow lines) shows the modified, nonequilibrium protocol used for the fast dynamical control of the three parameters $\alpha(t)$, $\beta(t)$, and $F(t)$ (followed by a reinitialization of the nonequilibrium state during time

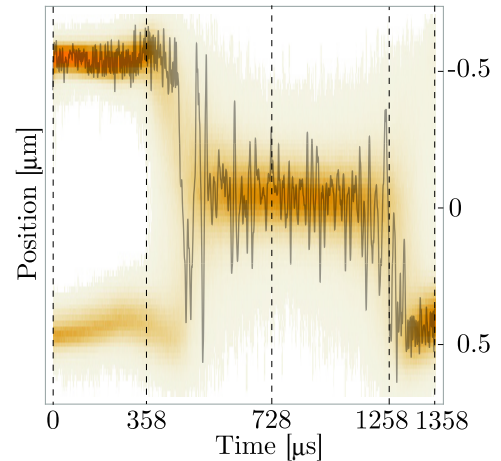


FIG. 3. Position distribution during erasure. Exemplary single stochastic trajectory of the nanoparticle (black line) from state “1” to state “0” measured during a nonequilibrium erasure protocol for $\epsilon = 4$. The orange cloud represents the corresponding position density for about 20 000 repetitions of the protocol. The protocol lasts $t_{\text{er}} = 1358\ \mu\text{s}$. In addition, $t_0 = 358\ \mu\text{s}$ marks the start of the usual equilibrium erasure, and $t_1 = 728\ \mu\text{s}$ identifies the introduction of the tilt, which reaches the maximum magnitude F_{max} at $t_2 = 1258\ \mu\text{s}$ [see also Fig. 2(b)].

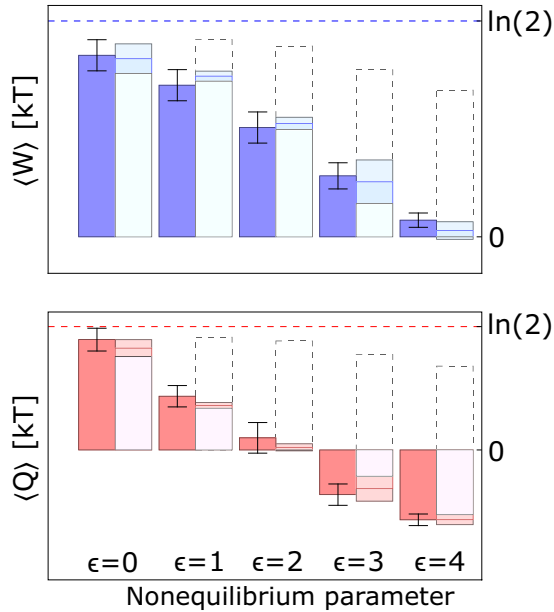


FIG. 4. Work and heat for the erasure of an initial, asymmetric nonequilibrium memory state. Average consumed work (blue bars) and dissipated heat (red bars) for various discrete values of the nonequilibrium parameter ϵ . As ϵ is increased, we observe a significant suppression of both quantities, below their corresponding equilibrium values for the same configuration (dashed white bars), in agreement with the generalized Landauer bounds (1) and (2) (solid white bars). The shaded areas represent the uncertainty on this value introduced by the experimentally determined asymmetry of the corresponding double-well potentials. The dashed horizontal lines indicate the $kT \ln 2$ limit valid for equilibrium symmetric memories. Error bars are 1- σ confidence intervals obtained for approximately 20 000 repetitions.

$\tau_{\text{rein}} = 1250 \mu\text{s}$). The modified protocol includes an initial fast modulation of the trap frequency and thus the potential shape. Such “decompression” phase allows one to pump heat from the environment to the system (and therefore suppress the dissipated heat) and to reduce the performed work [24]. The remaining “erasure” phase is essentially a standard equilibrium erasure.

Figure 3 displays the measured position distribution of the particle (yellow) over time for $\epsilon = 4$, along with one exemplary stochastic trajectory (black) showing a reset from bit “1” ($x < 0$) to bit “0” ($x > 0$). The initial asymmetry, in this case, is $p_{\epsilon=4}^0(0) = 56.7\%$ ($p_{\epsilon=4}^1(0) = 43.3\%$) for the potential well corresponding to state “0” (“1”). We set the memory to state “0” in time $t_{\text{er}} = 1358 \mu\text{s}$ with a success probability of $p_{\epsilon=4}^1(t_{\text{er}}) = 97.9\%$.

Figure 4 represents the average consumed work $\langle W \rangle$ (blue bars) and dissipated heat $\langle Q \rangle$ (red bars) for different values of the nonequilibrium parameter ϵ . We first observe that work and heat are smaller than the corresponding equilibrium values (dashed bars), taking the asymmetry of each initial state into account. We further clearly see that both $\langle W \rangle$ and $\langle Q \rangle$ decrease with increasing entropic distance from equilibrium of the initial memory state, in full agreement with Eqs. (1) and (2) (white bars): The consumed work is significantly reduced to $\langle W \rangle_{\epsilon=4} = (0.054 \pm 0.023) kT$, while the dissipated heat

even becomes negative for $\epsilon = 3$, $\epsilon = 4$, reaching $\langle Q \rangle_{\epsilon=4} = (-0.328 \pm 0.032) kT$, meaning that heat is absorbed (instead of dissipated) during information erasure for initial memory states that are sufficiently far from equilibrium. It is worth highlighting that both $\langle W \rangle$ and $\langle Q \rangle$ are close to their respective lower bounds given by Eqs. (1) and (2) (Table I), indicating that the reset protocol is nearly optimal. We also note that the effect of the initial nonequilibrium memory state on the energetics of the erasure process (difference between blue/red bars and dashed bars) is much larger than that coming from the asymmetry of the double-well potential (difference between blue/red dotted lines and dashed bars) (see also the discussion in Appendix A). Interestingly, the modified nonequilibrium reset protocol is close to optimality even in the case of non-ideal, asymmetric memories.

We additionally emphasize that our setup combines the ability to perform fast switches with the control over the equilibration time with the thermal environment, allowing for the implementation of a fast erasure protocol and at the same time the creation of NIM states. This highlights the importance of underdamped switches for future low-dissipation digital devices [49–52].

V. DISCUSSION AND OUTLOOK

In conclusion, we experimentally demonstrate that nonequilibrium memory states enable information erasure below the equilibrium Landauer limit for both work consumption and heat dissipation. This result is consistent with a generalized nonequilibrium erasure principle. Such reduction is physically possible when entropy and energy of the initial memory state are properly harnessed by the erasure protocol. Clearly, there is no way to avoid the overall thermodynamic cost during memory erasure as enforced by the second law for a full cycle. Yet, the existence of distinct preparation and erasure stages for nonequilibrium memories offers remarkable thermodynamic flexibility. The cost of erasure can thus be shifted to the preparation phase, as we have shown here. This property offers means to control the thermodynamics of logically irreversible operations [53] and manage heat production, for example, in proposed computer architectures where preparation and processing zones are spatially separated [54,55]. They may hence be implemented locally at no energetic cost and without heat dissipation or even entail cooling of the local environment. These results were made possible by the development of an adaptable optical micromanipulation scheme of a nanosphere trapped in a nonlinear potential and levitated in mild vacuum. This is an essential step toward bringing the same level of virtuous control known from optical micromanipulation of overdamped colloids in liquid [56] to levitated mesoscopic underdamped systems in vacuum. Optical levitation with dynamically shaped potentials is not only a versatile platform to investigate information thermodynamics but, together with the preparation of pure quantum states of nanoparticle motion [57–60], may also become a powerful tool for quantum state engineering beyond Gaussian states [43], and thus an approach to matter-wave interferometry in macroscopic quantum physics.

TABLE I. Energetics of the erasure of an asymmetric nonequilibrium initial memory state. Difference of relative entropies ΔI during the reset process and preparation energy $\Delta\langle U_{\text{pre}} \rangle = -\Delta\langle U_{\text{res}} \rangle$ of the nonequilibrium initial memory state, as well as work $\langle W \rangle$ and heat $\langle Q \rangle$ that enter the generalized Landauer bounds, Eqs. (1) and (2), for the data of Fig. 4 (all quantities are in units of kT).

ϵ	0	1	2	3	4
ΔI (kT)	0.572 ± 0.041	0.49 ± 0.05	0.363 ± 0.042	0.240 ± 0.033	0.042 ± 0.010
ΔU_{pre} (kT)	0.04 ± 0.08	-0.19 ± 0.08	-0.3 ± 0.1	-0.41 ± 0.07	-0.382 ± 0.039
$\langle W \rangle$ (kT)	0.583 ± 0.045	0.487 ± 0.048	0.351 ± 0.046	0.196 ± 0.042	0.054 ± 0.023
$\langle Q \rangle$ (kT)	0.620 ± 0.064	0.30 ± 0.06	0.024 ± 0.086	-0.22 ± 0.06	-0.328 ± 0.032

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DATA AVAILABILITY

The raw data for this study are available in the Phaidra Repository [61]. The scripts used to analyse the data are available upon reasonable request.

APPENDIX A: GENERALIZED LANDAUER BOUND

We reformulate Landauer's principle in terms of relative entropies between distributions to take into account the nonequilibrium nature of the initial state. Starting from the first and second law of thermodynamics [22], we can rewrite the mean work as

$$\langle W \rangle \geq \Delta\langle U \rangle - T \Delta S \equiv W_L, \quad (\text{A1})$$

where $\Delta\langle U \rangle$ is the change in internal energy and ΔS is the change in entropy. The time-dependent equilibrium Hamiltonian can be rewritten using the equilibrium distribution $H_{\text{eq}} = -kT \ln Z_{\text{eq}} - kT \ln \rho_{\text{eq}}$. Therefore, as shown in Ref. [24]

$$\begin{aligned} W_L &= \langle H_{\text{eq}} \rangle_i - \langle H_{\text{eq}} \rangle_f + T(S_f - S_i) \\ &= kT(\langle \ln \rho_{\text{eq}} \rangle_i - \langle \ln \rho_{\text{eq}} \rangle_f + \langle \ln \rho_f \rangle_f - \langle \ln \rho_i \rangle_i) \\ &= T(I(\rho_f, \rho_{\text{eq}}) - I(\rho_i, \rho_{\text{eq}})) = T \Delta I, \end{aligned} \quad (\text{A2})$$

with the relative entropy being defined as $I(\rho_{f/i}, \rho_{\text{eq}}) = k \langle \ln[\rho_{f/i}/\rho_{\text{eq}}] \rangle = k \int \rho_{f/i} \ln[\rho_{f/i}/\rho_{\text{eq}}] dx dv$. The mean heat is then given by [24]

$$\langle Q \rangle \geq W_L - \Delta\langle U \rangle \equiv Q_L, \quad (\text{A3})$$

using the sign convention $\Delta\langle U \rangle = \langle W \rangle - \langle Q \rangle$.

1. Asymmetry corrections

Let us next derive Landauer's bound by taking into account asymmetry in the potential and imperfect erasure. To that end, we approximate the double-well potential with a high-barrier by fitting both wells with harmonic oscillator potentials; we also ignore any potential momentum distribution

differences, i.e., we assume that the momentum distribution stays approximately thermal, for the theoretical bound. There are three relevant distributions: the equilibrium distribution ρ_{eq} , which is used to determine the equilibrium potential, the initial distribution ρ_i , and the final one ρ_f . Here, $p^{0/1}$ and $\sigma^{0/1}$ are distribution parameters that represent the amplitude and the width of the probability distribution, respectively. The position distribution for the equilibrium distribution is given by

$$\rho_{\text{eq}} = \begin{cases} p_{\text{eq}}^0 e^{-\beta \sigma_{\text{eq}}^0 (x-x_0)^2} / Z_{\text{eq}}^0, & \text{for } x > 0, \\ p_{\text{eq}}^1 e^{-\beta \sigma_{\text{eq}}^1 (x+x_1)^2} / Z_{\text{eq}}^1, & \text{for } x < 0. \end{cases} \quad (\text{A4})$$

The distribution parameters $p_{\text{eq}}^{0/1}, \sigma_{\text{eq}}^{0/1}$ are replaced by $p_i^{0/1}, \sigma_i^{0/1}$ when describing the initial out-of-equilibrium distribution, and by $p_f^{0/1}, \sigma_f^{0/1}$ for the final one. Here, we assume for simplicity that the position of the minimum does not change before/after the erasure protocol. The probability $p^{0/1}$ represents the probability of the particle being on either side of the wells, with $p^0 = 1 - p^1$. In order to recover the well-known equilibrium Landauer bound of $kT \ln 2$, $p_{\text{eq}}^0 = p_i^0 = 1/2$ and $p_f^1 = 1$. This treatment allows to recover the bounds accounting for asymmetry in the potential and a nonperfect erasure. We do assume, however, that $p_{\text{eq}}^{0/1} = p_i^{0/1}$: We improve the erasure by solely altering the entropy within the wells.

In thermal equilibrium, we can reconstruct the potential by using $V(x) = -kT \ln(p(x))$. Calculating explicitly the quantities in Eq. (A3), the energies are given by

$$\langle H_{\text{eq}} \rangle_i = \frac{kT}{2} \left(p_i^0 \frac{\sigma_{\text{eq}}^0}{\sigma_i^0} + p_i^1 \frac{\sigma_{\text{eq}}^1}{\sigma_i^1} \right) - kT(p_i^0 \ln p_{\text{eq}}^0 + p_i^1 \ln p_{\text{eq}}^1), \quad (\text{A5})$$

for the initial case and replacing $p_f^{0/1}$ for $p_i^{0/1}$ as well as $\sigma_f^{0/1}$ for $\sigma_i^{0/1}$ for the final case. The entropy difference is moreover

$$\begin{aligned} \Delta S &= -k \left(\int \rho_f \ln \rho_f dx - \int \rho_i \ln \rho_i dx \right) \\ &= k(p_i^0 \ln p_i^0 + p_i^1 \ln p_i^1 - p_f^0 \ln p_f^0 - p_f^1 \ln p_f^1) \\ &\quad + p_f^0 \ln Z_f^0 + p_f^1 \ln Z_f^1 - p_i^0 \ln Z_i^0 - p_i^1 \ln Z_i^1 \\ &= k \ln \frac{(p_i^0)^{p_i^0} (p_i^1)^{p_i^1}}{(p_f^0)^{p_f^0} (p_f^1)^{p_f^1}} + k \ln \frac{(Z_f^0)^{p_f^0} (Z_f^1)^{p_f^1}}{(Z_i^0)^{p_i^0} (Z_i^1)^{p_i^1}}. \end{aligned}$$

Using Eq. (A3), the expression for the generalized Landauer bound for our potential is thus

$$\begin{aligned} \frac{W_L}{kT} &= \frac{1}{2} \left(\frac{\sigma_{\text{eq}}^0 (p_i^0 \sigma_i^0 - p_i^0 \sigma_f^0)}{\sigma_i^0 \sigma_f^0} + \frac{\sigma_{\text{eq}}^1 (p_i^1 \sigma_i^1 - p_i^1 \sigma_f^1)}{\sigma_i^1 \sigma_f^1} \right) + \ln \frac{(p_{\text{eq}}^0)^{p_i^0} (p_{\text{eq}}^1)^{p_i^1}}{(p_{\text{eq}}^0)^{p_f^0} (p_{\text{eq}}^1)^{p_f^1}} - \ln \frac{(p_i^0)^{p_i^0} (p_i^1)^{p_i^1}}{(p_f^0)^{p_f^0} (p_f^1)^{p_f^1}} - \ln \frac{(Z_f^0)^{p_f^0} (Z_f^1)^{p_f^1}}{(Z_i^0)^{p_i^0} (Z_i^1)^{p_i^1}} \\ &= \frac{1}{2} \left(\frac{\sigma_{\text{eq}}^0 (p_f^0 \sigma_i^0 - p_i^0 \sigma_f^0)}{\sigma_i^0 \sigma_f^0} + \frac{\sigma_{\text{eq}}^1 (p_f^1 \sigma_i^1 - p_i^1 \sigma_f^1)}{\sigma_i^1 \sigma_f^1} \right) + \ln \frac{(p_f^0)^{p_f^0} (p_f^1)^{p_f^1}}{(p_{\text{eq}}^0)^{p_i^0} (p_{\text{eq}}^1)^{p_i^1}} + \frac{1}{2} \ln \frac{(\sigma_f^1)^{p_f^1} (\sigma_f^0)^{p_f^0}}{(\sigma_i^0)^{p_i^0} (\sigma_i^1)^{p_i^1}}, \end{aligned} \quad (\text{A6})$$

where in the last line we assumed $p_{\text{eq}}^0 = p_i^0$. There are three contributions to the work cost. The latter two are comprised of the logical component given by the probabilities and the internal potential width components. The first term is the energy difference between the initial locally squeezed distribution and the equilibrium one. The thermal energy $kT/2$ of the thermal distribution is reduced by the initial distribution in the equilibrium potential. The first fast step is thus accompanied by an energetic cost to work that makes the work larger than heat in the Landauer bound (up to $kT/2$). If we ignore the fast step, then this term would vanish, and work/heat would

be equal in the symmetric case and only differ because of the potential offsets. As the first step is faster than the movement of the particle, this work offset is just added numerically when retrieving the bound.

We further explicitly distinguish the contribution to the generalized Landauer's bound coming from the nonequilibrium character of the initial state (W_{NIMS}) from the one coming from the asymmetry of the potential (W_{asym}). To this end, we rewrite Eq. (A1) using $p^1 = 1 - p^0$, $\sigma^0 = \sigma^1$, $\sigma_i = r\sigma_{\text{eq}}$, with r the asymmetry parameter. We accordingly obtain

$$W_L/kT = \frac{1}{2} \left[2 \ln ((1 - p_f^0)^{1-p_f^0} (p_f^0)^{p_f^0} (1 - p_{\text{eq}}^0)^{p_f^0-1} (p_{\text{eq}}^0)^{p_f^0}) - \frac{1}{r} + \ln \left(\frac{1}{r} \right) + 1 \right] = \ln 2 - (W_{\text{NIMS}} + W_{\text{asym}})/kT. \quad (\text{A7})$$

In Table II, we report the experimental probabilities, as well as the asymmetry correction to the bound given in Eq. (A1). The probability p^0 (p^1) is calculated as the frequency of the particle being in well “0” (“1”) over the total number of repetitions. Errors on these parameters are evaluated by dividing the whole number of protocol repetitions into $n = 50$ statistically independent subsets from which the standard deviation is calculated and propagated.

APPENDIX B: TRAPPING SCHEME

We create the trap with a CFI TU Plan Fluor EPI 50x, NA = 0.8, WD = 1 mm (Nikon Corp.) objective in the vacuum chamber. Between the objective and the collimation lens, two razor blades, spaced by $\approx 100 \mu\text{m}$, are mounted to create the necessary linear unbalancing force. The electrode holder is a three-dimensional (3D)-printed, custom-designed element,

made of Accura Xtreme White (Proto Labs GmbH). The spacing is small enough that the electrodes are fixed in their position. The holder is screwed onto an MX25 translation stage that is controlled by a CU30 controller (Mechanics AG.). This allows for a three-dimensional movement of the electrodes with 2 mm total travel distance in each direction (see Fig. 5). Particles are trapped with the standard method using an ultrasonic nebulizer. We perform the experiment at a pressure in the vacuum chamber of $P = (30.0 \pm 0.3) \text{ mbar}$. The TEM₁₀ mode is created by means of a $512 \times 512 \text{ SLM}$ (Meadowlark, Inc.), which has been previously calibrated to account for the nonlinear response of the pixels and nonflatness of the active surface. A simple splitting of the SLM active area to create a π phase shift at the center of the beam (see Fig. 6) creates the wanted potential.

TABLE II. Experimental parameters. Probability p_0^{eq} of the particle starting in well 0 and probability p_0^f of the particle being in well 0 after the protocol, as well as correction $\langle W_{\text{asym}} \rangle$ to Landauer's bound given by the asymmetry of the potential as a function of the nonequilibrium parameter ϵ .

ϵ	p_0^{eq}	p_0^f	$\langle W_{\text{asym}} \rangle (\text{kT})$
0	0.634 ± 0.006	0.040 ± 0.001	0.121 ± 0.047
1	0.415 ± 0.006	0.004 ± 0.001	0.060 ± 0.016
2	0.423 ± 0.007	0.0046 ± 0.0014	0.081 ± 0.019
3	0.637 ± 0.021	0.033 ± 0.010	0.156 ± 0.070
4	0.566 ± 0.015	0.0210 ± 0.0035	0.222 ± 0.028

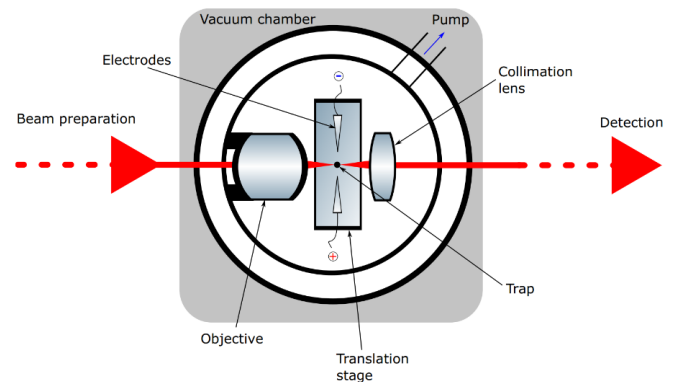


FIG. 5. Schematic representation of the vacuum chamber.

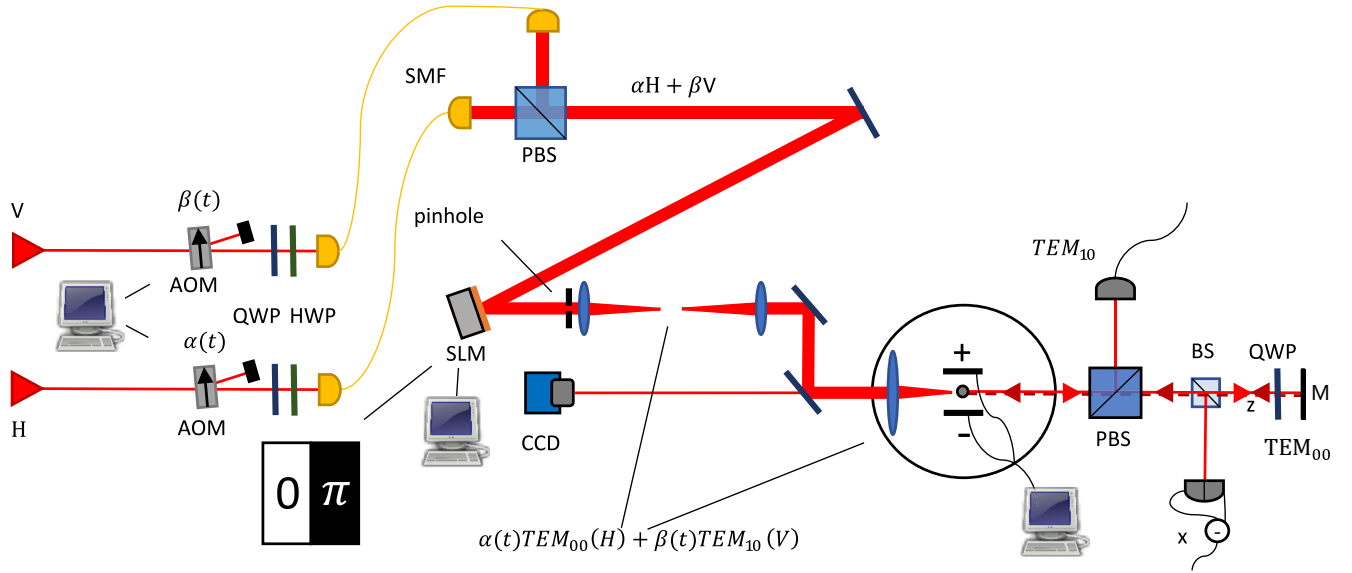


FIG. 6. Experimental setup: Two independent, orthogonal polarized laser beams [H: Mephisto (Coherent) + Keopsys Fiber Amplifier, V: Opus (Laser Quantum)] are modulated in intensity using independently controlled AOMs working in first-order diffraction and frequency shifted by 80 MHz. Light is coupled into the SMF and compensated in polarization, then expanded to a beam diameter of approximately 7 mm and recombined into a PBS. An SLM is programmed to create a TEM_{10} mode for the vertically polarized light. In the vacuum chamber, a microscope objective (NA = 0.8) creates the optical trap. Electrodes are placed around the focal point and can be positioned through a translation stage. The light from the trap is then collected and filtered in polarization. Vertically polarized light is used to monitor the TEM_{10} power, horizontally polarized light is split: 1% of the light is used to measure the particle position using a split-mirror detection scheme, the remaining 99% of the light is reflected into the vacuum chamber to create the standing wave through mirror M. List of acronyms: PBS, polarizing beam splitter; AOM, acousto-optical modulator; SLM, spatial light modulator; SMF, single-mode fiber; CCD, charged-coupled device fast camera; BS, beam splitter; HWP, half-wave plate; QWP, quarter-wave plate; M, mirror.

APPENDIX C: CALIBRATION

1. Position detection and detector calibration

The particle position is measured using a standard split-mirror detection scheme: The forward propagating TEM_{00} light mode is separated from the TEM_{10} mode, attenuated, and split along the vertical axis by means of a sharp-edged D-shaped mirror. Light is attenuated and sent to a pair of balanced InGaAs detectors (PDB425C; Thorlabs Inc.), where the difference signal ϕ_{det} is measured. The displacement of the particle is proportional to that signal when the particle is trapped close to the center of the potential ($\Delta x = c_{\text{center}} \phi_{\text{det}}$). In this condition, the calibration factor c_{center} can be calculated for each trapped particle using an independent measurement, where the power of the TEM_{10} mode is set to zero. By assuming harmonic motion in the TEM_{00} , we use the equipartition theorem: $m\omega_0^2 c_{\text{center}}^2 \sigma_{\phi_{\text{det}}}(\omega_0)^2 = kT$. Here, $m = (2.83 \pm 0.03) \times 10^{-18}$ kg is the mass of the silica nanosphere as provided by the manufacturer specifications, k is the Boltzmann constant, $T = 300$ K is the room temperature, and $\omega_0 = 2\pi(210.1 \pm 0.8)$ kHz is the particle frequency, for a power in the optical trap of $P_0 = (100 \pm 1)$ mW, where $\sigma_{\phi_{\text{det}}}(\omega_0) = (0.042 \pm 0.001)$ V is the standard deviation of the detector signal. We also measure a damping coefficient $\Gamma = 2\pi(32.0 \pm 0.3)$ kHz, leading to a quality factor of the oscillator of around 6.5. In Fig. 7, we show the linear response of the detector (the amplitude of the particle's signal with a standard deviation $\sigma_{\phi_{\text{det}}}(\omega_0)$), as a function of the overall power in the detector (measured in volts). The linear dependence shows

that we can renormalize the particle's signal to the power during the protocol. For this measurement, the particle is trapped at a pressure of 0.1 mbar and only in the TEM_{00} mode at constant power ($\alpha = 1, \beta = 0, F = 0$). For this configuration, we obtain

$$c_{\text{center}} = \sqrt{\frac{kT}{m}} \frac{1}{\sigma_{\phi_{\text{det}}} \omega_0} = (6.90 \pm 0.17) \times 10^{-7} \text{ m/V}.$$

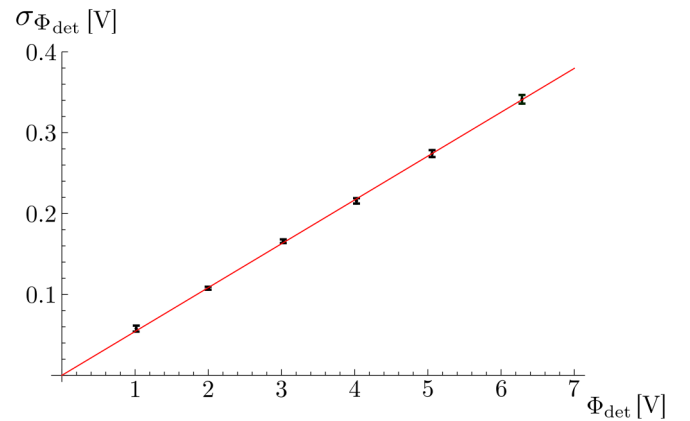


FIG. 7. Power dependence of detector sensitivity. $\sigma_{\phi_{\text{det}}}$ as a function of the power in the detector shows a linear dependence, verifying the operation of the detector in the linear response regime. The optical power is attenuated right in front of the detector, for a fixed trap frequency of ω_0 .

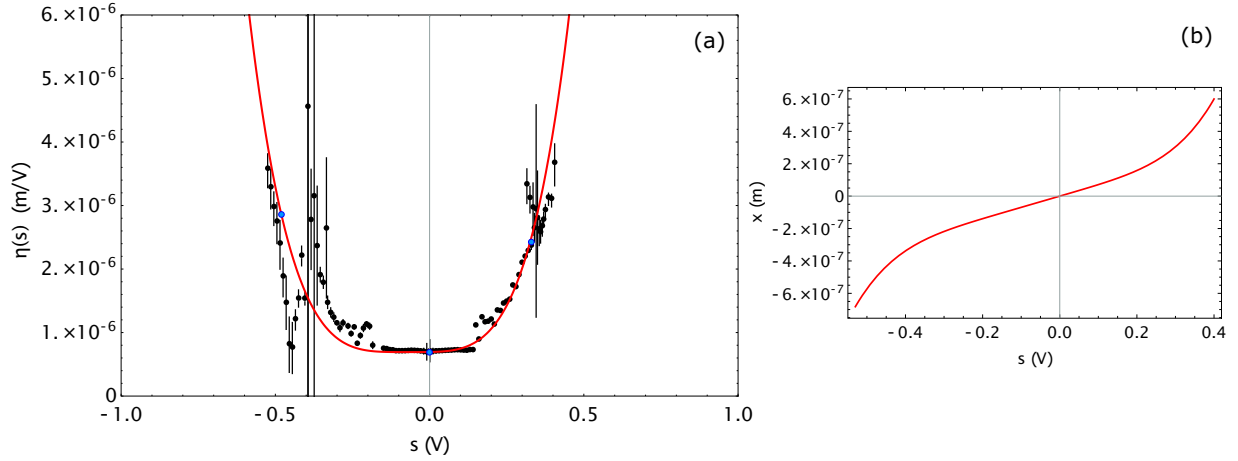


FIG. 8. Detector sensitivity. (a) Differential calibration factor as a function of the detector response. Black points: experimental values. Error bars are $1\text{-}\sigma$, evaluated with a Montecarlo simulation. Blue points: calibration factor evaluated using the equipartition theorem in approximately harmonic traps, with error bars smaller than the point size. Red: fit of all data points with a quartic function. (b) Correspondingly calibrated position as a function of the detector response (uncalibrated position).

However, when the particle is significantly off-centered (for example, when the particle is in either of the two wells of the TEM_{10} mode), the sensitivity (i.e., the calibration factor) of the detection scheme is reduced, because less light from the TEM_{00} mode is scattered by the particle. In fact, if we repeat the above calibration for when the particle is in the left and right wells, we get

$$\begin{aligned} c_{\text{left}} &= (2.95 \pm 0.11) \times 10^{-6} \text{ m/V}, \\ c_{\text{right}} &= (2.62 \pm 0.10) \times 10^{-6} \text{ m/V}. \end{aligned}$$

Thus, we investigate the dependence of the calibration factor on the particle position $c(x)$ for the relevant spatial interval. This is possible by analyzing the particle dynamics in an equilibrium scenario, as discussed below. The data used to calibrate the effect stems from the reinitialization part of the protocols, i.e., is taken independently of the erasure protocol and with precisely the same experimental configuration and particle. During the reinitialization phase, we obtain trajectories with the particle equilibrated in both wells. We obtain this data for two different powers, with the potential set to “high-power” and “low-power.” We label $s(t)$ the uncalibrated particle trajectory registered by the detector. The goal is to reconstruct the calibrated particle trajectory in real space $x(t)$. We define $\eta(s) = dx/ds$ as the differential calibration factor, which we expect to correspond to the calibration factor c_i as discussed above ($i = \text{left, right, center}$). Accordingly, $\partial V/\partial s = \partial V/\partial x \times \partial x/\partial s = F(x)\eta(s)$. Let us first consider an isolated particle (no friction, no force noise) moving in a potential landscape and detected in a scheme with a position-dependent calibration factor:

$$F(x) = \frac{1}{\eta(s)} \times \frac{\partial V}{\partial s}, \quad (\text{C1})$$

$$F(x)/m = \frac{\partial}{\partial t}(\dot{x}) = \frac{\partial}{\partial t}(\eta(s)\dot{s}) = \frac{\partial \eta}{\partial s}\dot{s}^2 + \eta(s)\ddot{s}, \quad (\text{C2})$$

$$= \frac{1}{\eta(s)} \left(\frac{\partial \eta^2}{\partial s} \times \frac{\dot{s}^2}{2} + \eta^2 \ddot{s} \right). \quad (\text{C3})$$

From this, we can deduce how Newton’s laws are transformed for the uncalibrated trajectories in dependence on the differential calibration factor:

$$\frac{\partial V}{\partial s} = m \frac{\partial \eta(s)^2}{\partial s} \times \frac{\dot{s}^2}{2} + m \eta^2(s) \ddot{s}. \quad (\text{C4})$$

Clearly, the detected uncalibrated trajectory allows for determining the calibration factor using the kinetic quantities in this equation. Consider a particle weakly coupled to and in equilibrium with an environment. For a sufficiently large ensemble of data where a particle crosses a specific position, the above equations hold on average, as random force noise and friction cancel on average. Accordingly,

$$\left\langle \frac{\partial V}{\partial s} \right\rangle = \frac{\partial \eta(s)^2}{\partial s} \times \left\langle \frac{\dot{s}^2}{2} \right\rangle + m \eta^2(s) \langle \ddot{s} \rangle. \quad (\text{C5})$$

The calibration factor is now expressed in terms of a first-order differential equation with uncalibrated quantities, which are known: the reconstructed potential derivative, the average kinetic energy, and the average acceleration when the particle is in the TEM_{00} mode and in the wells of the TEM_{10} high- and low-“power” mode. We numerically solve the differential equation to reconstruct $\sqrt{\eta^2(s)}$.

In Fig. 8, we show the differential calibration factor as a function of the voltage response of the detector. The data are obtained by stitching the numerical results evaluated in various conditions of the potential: When the particle is in the two wells, both at “high-power” (at $\epsilon = 4$) and at “low-power” (at $\epsilon = 0$), as well as when the particle is in the TEM_{00} potential. While this allows for the full relevant range of data to be covered, it results in two apparent discontinuities and corresponding large error bars due to numerical mismatch in the different reconstructions. In fact, for each component, we manually adjust the position offset ($V(s + \xi) = \xi_0$) to minimize discontinuities. We evaluate error bars on the points by performing a Montecarlo simulation where we randomly generate ξ in an interval around ξ_0 . To compare this dynamics-based full-range calibration method with the standard measurement for harmonic potentials described before,

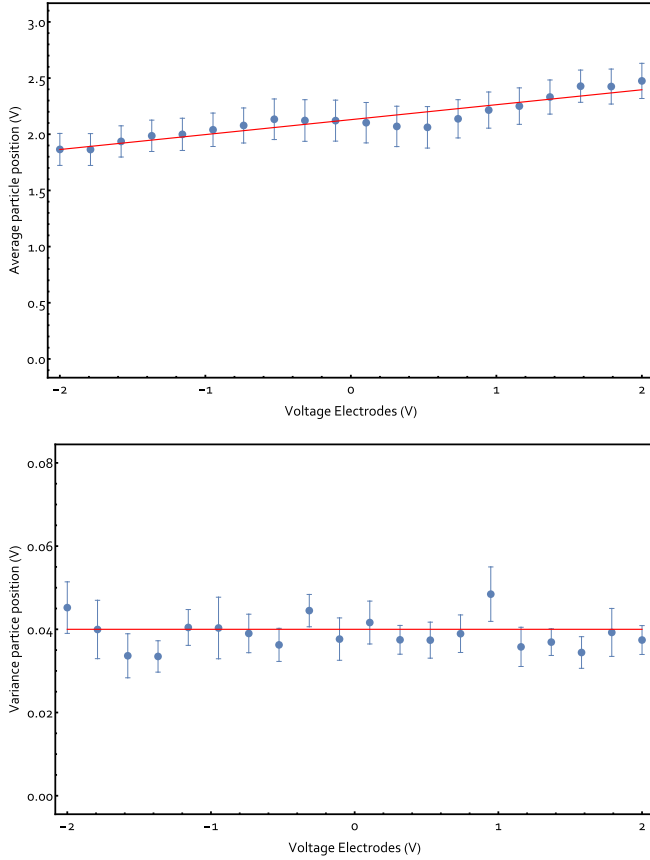


FIG. 9. Calibration of the electrodes. (Top) Average of the center-of-mass mean position of the particle as a function of the voltage applied to the electrodes. (Bottom) Average of the particle motion variance, as a function of the voltage applied to the electrodes. Data points and error bars are obtained over ten repetitions of acquiring 10 ms time traces with varying electrode voltage. In red, the linear fit, obtaining the slope $c = (0.133 \pm 0.0028) V_{\text{detector}}/V_{\text{electrodes}}$ (top) and $c = (0.000 \pm 0.028)$ (bottom).

Fig. 8 also features the values c_{left} and c_{right} , and c_{center} as blue dots showing a good agreement between the two calibration methods.

We interpolate the data points with the function $\alpha_4(s - s_0)^4 + \alpha_0$, which preserves the expected flatness of the calibration function around zero. In Fig. 8(b), we show the integrated calibration factor $\eta_0(x) = \int_{s_{\min}}^x \eta(s') ds'$, which allows us to determine the real space trajectory for any of our measurements in postprocessing. Every measured timetrace is rescaled to account for power fluctuations, which are monitored during the whole measurement, and their linear effect on the calibration in Fig. 7. Specifically, the data are normalized to a power P_0 corresponding to a measured voltage of $\phi_{\text{det,max}} = (5.51 \pm 0.02) \text{ V}$.

2. Electrode calibration

We calibrate the displacement force $F = c_f \phi_f$ due to the voltage ϕ_f applied to the electrodes by measuring the displacement of the particle in the TEM₀₀ mode as a function of the voltage applied to the electrodes. The electrode's voltage is increased from $\phi_f = -2.0 \text{ V}$ to $\phi_f = 2.0 \text{ V}$ in multiple

TABLE III. Experimental parameters. Probability of finding the particle in the left well at the beginning of the protocol (p_0^{eq}), and probability of finding the particle in the left well at the end of the protocol (p_f^0) with and without applying the external force of maximum magnitude F_{max} .

$F_{\text{max}} (\text{V})$	p_{eq}^0	p_f^0
0.86	0.566	0.021
0	0.572	0.535

steps using a voltage amplifier. At every step, we measure the mean position of the particle in the tweezer. A repetition of 10 times decreases the statistical error and compensates for position drifts. Afterward, the measurement is fitted with a linear function (see Fig. 9) from which we extract the parameter c_f . We get

$$\frac{\langle \phi_{\text{det}} \rangle}{\phi_f} = \frac{\langle x \rangle / c_{\text{det}}}{F / c_f} = \frac{1}{m\omega_0^2} \times \frac{c_{\text{det}}}{c_f}.$$

Note that every particle has a different number of charges and thus needs a different calibration. All protocols presented here are taken for the same particle. The force calibration factor determined for this particle is $c_f = (5.1 \pm 0.3) \times 10^{-14} \text{ N/V}$. The maximal tilting force applied in the protocol is $F_{\text{max}} = 4.4 \times 10^{-14} \text{ N}$ and corresponds to $\phi_{f,\text{max}} = F_{\text{max}} / c_f = 0.86 \text{ V}$. Additionally, we tested multiple repetitions of the equilibrium protocol where the force was not applied. For example, we show in Table III, for the case of $\epsilon = 4$, that not applying the external force does not change significantly the probability distribution between the beginning and the end of the protocol.

APPENDIX D: POTENTIAL MODEL

The calculation of work applied during a protocol is based on the knowledge of the optical potential at every time. We model the full time dependence based on our time-dependent parameter settings and the reconstruction of the potential for relevant, specific parameter settings, corresponding to snapshots of the protocol. The potential reconstruction for these snapshots is obtained from a series of independent measurements: We first let the particle equilibrate to the respective potential landscape. Then, we measure the particle's position probability distribution $p(x)$, taking a timetrace, which is sufficiently long (100 ms) to use the ergodic assumption safely. The potential is then reconstructed by using: $V(x, t_0) = -kT \ln(p(x, t_0))$ with Boltzmann constant k and temperature T . We measure the potential for various powers in TEM₁₀ and fixed TEM₀₀ power. Measurements are then sensitivity corrected, calibrated (see Appendix C 1), and fitted to find a model for the potential depending on the power of TEM₀₀ and TEM₁₀ (see some examples of measured potentials in Fig. 10). Every potential is fitted with a function, which approximates the optical potential as a combination of TEM₀₀ and TEM₁₀ modes:

$$f(x) = c - \frac{8\alpha(x - x_0)^2 e^{-\frac{2(x-x_0)^2}{\sigma^2}}}{\pi\sigma^6} - \frac{2\beta e^{-\frac{2(x-x_1)^2}{\sigma^2}}}{\pi\sigma^4}. \quad (\text{D1})$$

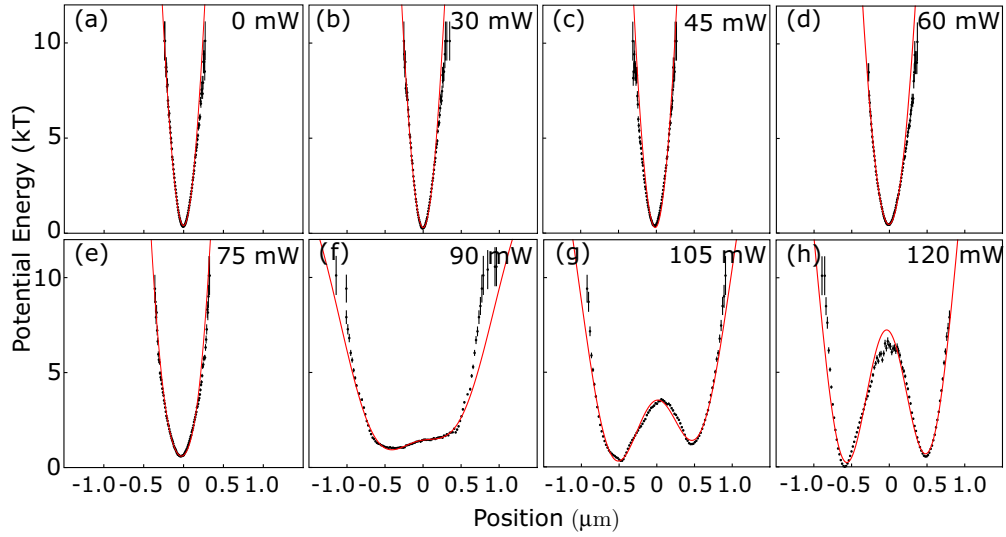


FIG. 10. Potential reconstruction. Reconstruction of $U(x)$, fixing the power of the TEM_{00} after the vacuum chamber to $\alpha_0 = 50$ mW, and varying the power of TEM_{10} from $\beta = 120$ mW [panel (h)] to $\beta = 0$ mW [panel (a)].

We first identify parameter β by fitting our potential when TEM_{10} is off [Fig. 10(a)]. We use a “jumping” potential (i.e., when the particle spontaneously jumps between the wells, Fig. 10(h)) to infer $|x_0 - x_1|$, which identifies the asymmetry in our model, which we assume constant for the whole experiment. We then fit the remaining potentials to find and interpolate parameters α , x_0 , σ , and c as a function of power in the mode TEM_{10} . This process allows us to program $V(x, t)$ for our protocol via the input parameters.

APPENDIX E: EXPERIMENTAL PROTOCOL AND INITIAL STATE PREPARATION

Every protocol timetrace is measured using a split detection scheme and sampled at 6.25 MS/s. A calibrated timetrace

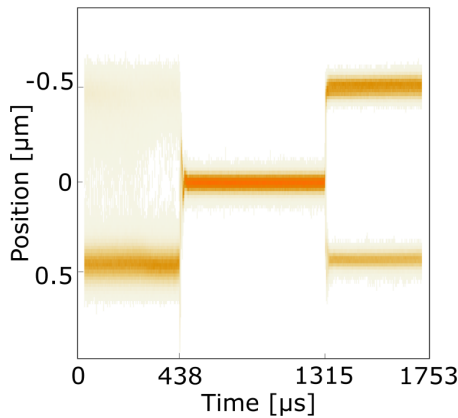


FIG. 11. Probability density of the reinitialization part of the protocol. Average position over $\approx 10\,000$ protocol repetitions. In the first part, the particle is in the “low” potential, after the erasure protocol. At $t = 438$ μs , the power of TEM_{00} is maximized. After $t = 1315$ μs , the TEM_{10} power is also maximized. At the end of the reinitialization, the particle is reequilibrated to a “high” power potential.

is identified by $x^l(t_i)$ (i is the i th time step, l is the l th repetition of the protocol). A full experimental protocol consists of three parts [Figs. 6(b) and 6(c)]: a preparation phase, the erasure protocol, and the reinitialization phase.

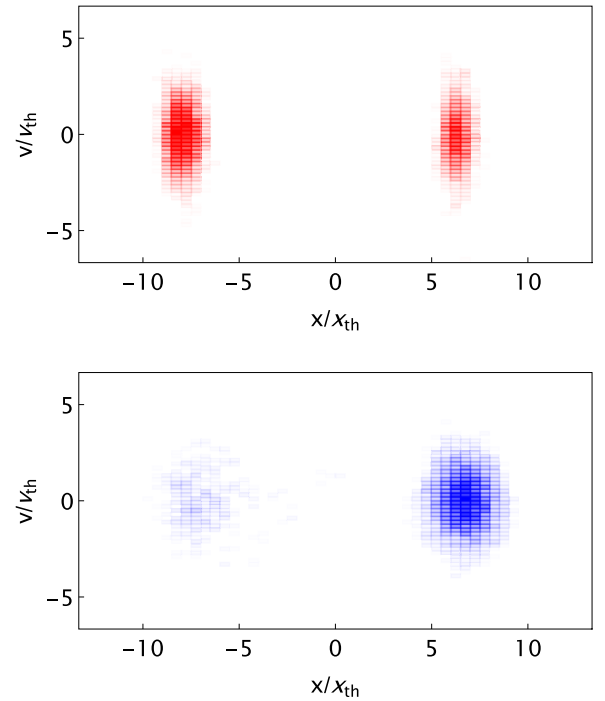


FIG. 12. Initial and final phase-space distributions. Top: Nonequilibrium initial memory state prepared at time $t = 0$ for $\epsilon = 4$. Bottom: Final equilibrium state after erasure obtained at time t_{er} . In both plots, position and velocity are normalized to $x_{\text{th}} = \sqrt{k_B T / (m(2\pi\nu_0)^2)}$ and $v_{\text{th}} = \sqrt{k_B T / m}$, where $\nu = 80$ kHz, $m = 2.83$ fg, and $T = 300$ K. The thermal squeezing present in the initial state is clearly observable in contrast to the final state.

TABLE IV. Relative entropy between the equilibrium state ρ_{eq} and $\rho(x, v, 0)$ and $\rho(x, v, t_{\text{er}})$, respectively, in dependence on the parameter ϵ . Equation (2) in the main text relates their difference ΔI to the expected work during the Landauer protocol $\langle W \rangle$.

ϵ	0	1	2	3	4
$I(0)(\text{kT})$	0	0.29 ± 0.03	0.39 ± 0.04	0.55 ± 0.06	0.53 ± 0.05
$I(t_{\text{er}})(\text{kT})$		0.77 ± 0.08	0.75 ± 0.08	0.79 ± 0.08	0.57 ± 0.06
$\Delta I(\text{kT})$		0.49 ± 0.05	0.36 ± 0.04	0.24 ± 0.03	0.04 ± 0.01

In the first part, we perform a fast ($t_{\text{up}} = 3 \mu\text{s}$) increase in the power of the trap, followed by a slower ($t_{\text{neq}} = 358 \mu\text{s}$) decrease back to the memory potential; the second part is the equilibrium erasure, where TEM_{10} is quasiadiabatically modified first to decrease and then increase again the potential barrier ($t_{\text{eq}} = 1358 \mu\text{s}$). At $t_1 = 728 \mu\text{s}$, we introduce a linear increase of the electrostatic force that reaches its maximum F_{max} at $t_2 = 1258 \mu\text{s}$, decreasing again linearly till the end of the protocol. The third part is a reset that restarts the particle in its original (nonequilibrium) distribution ρ_{neq} . This part is not evaluated for the work and heat consumption. First, the particle is allowed to equilibrate to the final position, then TEM_{00} is rapidly increased to its maximum value. Finally, TEM_{10} is also increased rapidly. This sets the particle to either of the two wells with a probability that depends on the symmetry of the potential. Finally, the particle is allowed to equilibrate to this potential. We show in Fig. 11 the probability density during

the reinitialization part of the protocol for the $\epsilon = 4$ dataset. Overall, the entire dataset lasts $t_{\text{tot}} = 3116 \mu\text{s}$. The double-well potential is consistently set to the same power ratio between the two modes (TEM_{10} is 2.7 times that of the TEM_{00} measured right after the vacuum chamber). The parameter ϵ is set so that the trap power is $(\epsilon + 1)$ times stronger than the power for the equilibrium potential. In Fig. 12, we show the phase-space distribution of the motional state at the beginning (top) and at the end of the erasure protocol (bottom), reflecting the thermally squashed nature of the initial state.

APPENDIX F: QUANTIFYING THE DISTANCE FROM EQUILIBRIUM

We further quantify the deviation from equilibrium via the relative entropy of the initial and final states in Table IV. In order to extract $I(t) = -k \iint \rho(x, v, t) \ln \rho(x, v, t) / \rho_{\text{eq}}(x, v, t) dx dv$, we additionally assume that the only significant change to the phase space distribution in equilibrium is encoded in the spatial distribution $\rho(x, v, t) = p_v p(x, t) dx \times p(x, 0)$ is then reconstructed using the distribution of $x(t = 0)$, and $p(x, t_{\text{er}})$ is reconstructed using the ensemble distribution of $x(t = t_{\text{er}})$. Since we do not have direct access to ρ_{eq} during the protocol, we need to use a reference distribution coming from the protocol runs where the electrode force is not activated. The additional assumption is that the potential landscape (and thus the probability distribution) stays the same between these two experimental runs.

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