

## Spin-mechanical thermal machines

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We present a framework for developing quantum thermal machines—a quantum battery and a quantum Otto engine—using a single nitrogen-vacancy (NV) center coupled to a low-frequency mechanical oscillator. These machines are engineered to operate in a regime where the oscillator’s frequency is much smaller than the NV center’s spin transition and relaxation rates. By employing carefully timed microwave pulses and rapid optical spin resets, we enable on-demand, coherent energy transfer between the NV center and the oscillator. In the quantum battery mode, energy is incrementally stored in or extracted from the oscillator’s long-lived motional states via discrete spin flips. Furthermore, spin-driven protocols facilitate a cooling mechanism that reduces the oscillator’s effective phonon population below thermal equilibrium, even for frequencies lower than the spin’s decoherence scale, effectively engineering a cold bath. Integrating this spin-cooling mechanism into an Otto cycle while maintaining contact with a hot reservoir enables the realization of a single-bath quantum heat engine. This engine generates mechanical work with near-optimal efficiency in finite time without requiring environmental decoupling.

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## I. INTRODUCTION

The intersection of thermodynamics and quantum mechanics has opened fields of research for understanding energy transfer at the quantum level [1–10]. In classical thermodynamics [11], heat refers to a form of energy that, when transferred, produces entropy. This energy transfer generally occurs through thermal conduction, convection, or radiation, and is fundamental for maintaining thermal equilibrium in a system. On the other hand, work is defined as the transfer of energy that does not increase entropy, usually when the system moves or changes shape. Carnot demonstrated two centuries ago [12], in 1824, that work can be extracted via reversible cycles where a working fluid alternates contact between hot and cold reservoirs. Miniaturizing thermodynamic machines toward the quantum regime may have practical applications in energy extraction and storage, such as quantum batteries [13–17] and quantum heat engines (QHEs) [4,5,18–22]. Extracting work from quantum thermal engines has been studied in various quantum systems theoretically and experimentally [4,5,7,18–26]. These include single qubits coupled to thermal baths [5,19,27–30], quantum oscillators interacting with reservoirs [24,31–33], and systems where the statistics of identical bosons or fermions are manipulated [34–37]. Despite these

advancements, achieving total control over the interaction between quantum systems and their environments remains a significant challenge. Controlling energy flow at the quantum scale promises significant advancements in nanotechnology [7,23,25,26] and materials science [38,39], particularly in quantum thermometry [23,40–43]. Significant progress has been made in manipulating spin-mechanical systems for applications in quantum sensing [44,45], simulations [46,47], and energy transduction [48]. However, a large disparity typically exists between the bare motional and spin frequencies, which are trying to be addressed by dressing the spin transitions via driving fields [49,50].

In this work, we propose to construct a quantum battery and heat engine using a high- $Q$  diamond trapped in high vacuum [51–53]. The center-of-mass motion of the particle, with frequencies around 10–100 Hz, serves as the working fluid. Embedded within the diamond is a spin system operating at GHz frequencies. We engineer spin-mechanical coupling via magnetic-field gradients to achieve coherent exchange between the spin and the particle’s motion. Crucially, the spin can be reset through optical pulses—a spin polarization process that irreversibly reduces the spin’s entropy [54].

## II. SIGNIFICANCE

Typically, thermal machines require the working medium to be in alternate contact with a hot and a cold bath. Arranging for such time-dependent couplings can be very challenging in practice [55–57]. Our work is significant in that we replace the step of coupling the working medium to the cold bath by the illumination of the system with green light. Application of the green light polarizes the NV and reduces the entropy. By controlling the timing and durations of the illumination by green light we can provide detailed control of the removal of

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entropy from the motional system via spin-mechanical coupling.

There are many works already studying quantum batteries using spin-mechanical systems, but all of these focus on storing the energy in a spin ensemble [58–60]. Our work is significant in that we propose to store the energy in a bosonic motional degree of freedom, which can possess extremely long coherence times as compared with spin ensembles. Thus, our battery can have a much longer lifetime than traditional spin-based quantum batteries.

Our designs operate via spin-mechanical coupling. There have been a number of previous works that also make use of spin-mechanical coupling, but in most of those cases the spin transitions need to be nearly resonant with the mechanical frequencies involved. This can be extremely challenging to achieve in practice, as mechanical oscillators with high-quality factors in the MHz–GHz regime are challenging to manufacture [61]. Our work is significant in that we show a scheme to couple spins with transitions in the MHz–GHz frequencies to mechanics operating in the Hz regimes via a dressing of the spin transitions. This opens up a great flexibility in control, making the protocols viable for a range of mechanical frequencies ranging from Hz to MHz.

Although several schemes have been developed to cool the mechanics using spins, especially using NV centers, the cooling is typically only effective when the mechanical frequency is comparable to or greater than the dephasing rate of the spin. This means that the spin dynamics retains some coherence over several mechanical oscillator periods [49,62,63]. Our work is significant in that by the alternating application of microwave and laser pulses to the spin, we can demonstrate cooling of the mechanical oscillator even when the spin coherence time is far shorter than the oscillator period.

### III. THE MODEL

We consider a nanodiamond containing an NV center whose position is harmonically trapped along  $\vec{e}_z$  at a low frequency  $\omega_m$ . The quantum Hamiltonian describing the self-energy of the phononic mode of the center-of-mass motion of the diamond (particle moving in a harmonic oscillator potential) is given by

$$\mathcal{H}_m/\hbar = \omega_m a^\dagger a, \quad (1)$$

where  $a$  ( $a^\dagger$ ) are the destruction (creation) operators for the phonon. We apply a static external magnetic field of the form  $\vec{B} = (B_0 - G_z z)\vec{e}_z$  to the trapped diamond, where  $B_0$  is a homogeneous magnetic field and  $G_z$  is the value of the magnetic-field gradient along  $\vec{e}_z$  at the vertical position  $z$ . We further assume the NV axis to also lie along  $\vec{e}_z$ . The spin Hamiltonian of an NV in an external magnetic field is a combination of the so-called zero field splitting term and a Zeeman term [54]:

$$\mathcal{H}_{\text{spin}}/\hbar = D S_z^2 + \gamma \vec{B} \cdot \vec{S}, \quad (2)$$

where  $D$  is the zero-field splitting of the electronic energy levels of the NV center,  $\vec{S}$  the spin-1 vector operator with components ( $S_x, S_y, S_z$ ), and  $\gamma \sim 28 \text{ GHz/T}$  is the gyromagnetic ratio of the electron. Inserting the external magnetic field in

Eq. (2), from our assumed form given above, we obtain

$$\mathcal{H}_{\text{spin}}/\hbar = D S_z^2 + \gamma B_0 S_z - \gamma G_z z S_z, \quad (3)$$

where  $z$  initially appears as a classical  $c$ -number coordinate of the nanodiamond. However, since we are considering quantized motion as described by Eq. (1),  $z$  must be reexpressed in terms of the annihilation and creation operators as  $z = (a + a^\dagger) \times z_{\text{zpf}}/\sqrt{2}$ , where  $z_{\text{zpf}} = \sqrt{\hbar/(m\omega_m)}$  is the zero point fluctuation length and  $m$  is the mass of the nanodiamond. We observe that the right-hand term of Eq. (3) couples the spin and the mechanics. In the case of a high-magnetic field, the transition frequencies between the two allowed transitions,  $m_S : 0 \leftrightarrow \pm 1$ , are very different, and thus we can restrict ourselves to consider the effective two-level system (TLS):  $m_S = 0 \leftrightarrow -1$ . We identify the NV sublevels as  $|g\rangle \equiv |m_S=0\rangle$  and  $|e\rangle \equiv |m_S=-1\rangle$ , and define the Pauli operators  $\sigma_z = |g\rangle\langle g| - |e\rangle\langle e|$ ,  $\sigma_+ = |e\rangle\langle g|$ ,  $\sigma_- = |g\rangle\langle e|$ , so that  $|e\rangle\langle e| = (1 - \sigma_z)/2$ . With this convention, the bare TLS Hamiltonian reads  $H_{\text{TLS}}/\hbar = \frac{\omega_0}{2}(1 - \sigma_z)$  (up to an irrelevant constant), where  $\omega_0 = \gamma B_0 - D$ .

Without loss of generality, we restrict ourselves to  $B_0 < 1024 \text{ G}$  (the magnetic field where the  $m_S = 0$  and  $m_S = -1$  levels cross), to preserve  $|0\rangle$  as the lowest energy state.

We now include an additional microwave drive of the TLS with strength  $\Omega$  and at the frequency  $\omega_L$ , given by  $\mathcal{H}_D/\hbar = \Omega(\sigma_+ e^{-i\omega_L t} + \sigma_- e^{i\omega_L t})$ . Moving to the frame rotating at the microwave frequency  $\omega_L$ , generated by  $U(t) = \exp[-i\omega_L t(1 - \sigma_z)/2]$ , the free TLS term becomes  $\frac{\Delta}{2}(1 - \sigma_z)$  with detuning  $\Delta \equiv \omega_L - \omega_0$ . Under the rotating-wave approximation, the drive reduces to  $\Omega\sigma_x$ , yielding Eq. (4):

$$\mathcal{H}/\hbar = \omega_m a^\dagger a + \frac{\Delta}{2}(1 - \sigma_z) - \frac{g}{2}(a + a^\dagger)(1 - \sigma_z) + \Omega\sigma_x, \quad (4)$$

where the TLS Pauli operators in Eq. (4) are defined relative to the basis states  $|s = g\rangle \equiv |m_S = 0\rangle$  and  $|s = e\rangle \equiv |m_S = -1\rangle$ , and with the effective TLS frequency  $\omega_0 = \gamma B_0 - D$ , spin detuning  $\Delta = \omega_L - \omega_0$  and spin-mechanical coupling strength  $g = \gamma G_z z_{\text{zpf}}$ . For this kind of hybrid system,  $\omega_m$  lies in the Hz–MHz range,  $\Delta$  and  $\Omega$  can both range from 0 to GHz scale, and  $g$  can start from 0 and can reach a predicted value of the order 100 kHz [64–66].

We see that the third term in Eq. (4) vanishes when the effective TLS is in the  $|g\rangle$  state, as  $m_S = 0$  for this state, and the spin energy is independent of the magnetic field.

We assume that the mechanical degree of freedom has an equilibrium position  $\langle z \rangle = 0$  when the spin state  $|s = g\rangle$ . When  $|s = e\rangle$ , the mechanical system's equilibrium position is shifted to  $z_0 = \gamma G_z / m\omega_m^2$ , as depicted in Fig. 1(a). Therefore, in this hybrid system, we naturally obtain a spin-mechanical interaction, which produces a force (or not) on the diamond depending on the spin state of the NV center. We now consider modulating the strength of the microwave drive in time in a discrete binary pulsed on/off fashion, i.e.,  $\Omega(t) = \Omega_L f(t)$ , where  $\Omega_L$  is fixed in time and  $f(t) = \{0, 1\}$  is a binary function. By constructing different pulse trains  $f(t)$ , we will be able to engineer complex dynamics for the hybrid system.

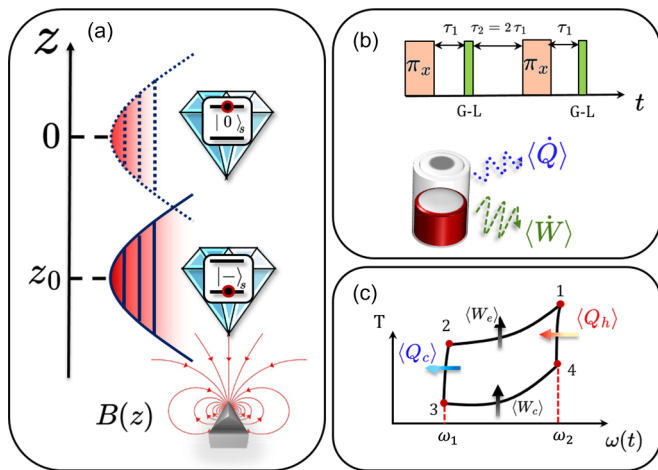


FIG. 1. (a) Proposed setup: A harmonically trapped diamond containing an NV center is subjected to a gradient magnetic field  $B(z) = B_0 - G_z z$ . In the high-magnetic-field limit, the NV center can be reduced to a spin-1/2 system with states  $|0\rangle_s$  and  $|-1\rangle_s$ . The spin-mechanical coupling  $\sigma_z(a + a^\dagger)$  induces a spin-dependent shift in the trap position, enabling the creation of two thermal machines illustrated in panels (b) and (c). (b) Quantum battery: Energy is stored in the high-quality mechanical oscillator due to the spin-dependent coupling and the long mechanical coherence time. The protocol alternates between transverse spin pulses, free evolution, and spin resetting via optical means, effectively transferring and storing energy in the mechanical mode. (c) Quantum heat engine: By adiabatically modulating the mechanical frequency between  $\omega_1$  and  $\omega_2$  and tuning the microwave drive on the spin, an effective cold reservoir is engineered. This allows energy extraction from the mechanical oscillator using a single thermal bath at ambient temperature.

We describe the evolution of the full spin-mechanical system with a density matrix  $\rho$ , including decay and dephasing of the TLS and damping of the motion via the master equation:

$$\frac{d}{dt}\rho = \frac{1}{i\hbar}[\mathcal{H}, \rho] + \sum_k \mathcal{L}_{x_k}[x_k]\rho + \sum_k \mathcal{L}_{y_k}[y_k]\rho, \quad (5)$$

$$\mathcal{L}_{x_k}[x_k]\rho = \frac{\gamma_{x_k}}{2}[(\bar{n}_{x_k} + 1)\mathcal{L}[x_k]\rho + \bar{n}_{x_k}\mathcal{L}[x_k^\dagger]\rho], \quad (6)$$

$$\mathcal{L}_{y_k}[y_k]\rho = \frac{\gamma_{y_k}}{2}\mathcal{L}[y_k]\rho, \quad (7)$$

where  $x_k = \{a, \sigma_-\}$ , with rates  $\gamma_{x_k} = \{\gamma_m, \gamma_1\}$ , describes the coupling of the motion and spin to their thermal baths, and  $y_k = \{\sigma_z, \sigma_-\}$ , with rates  $\gamma_{y_k} = \{\gamma_2, \Gamma_{GL}\}$ , describes spin dephasing and the process of spin polarization by the green light illumination. Here,  $\gamma_m$  is the mechanical damping rate giving the oscillator a mechanical quality factor of  $Q = \omega_m / (2\pi \times \gamma_m)$ ,  $\bar{n}_{x_k} = (\exp[\hbar\omega_k/k_B T_k] - 1)^{-1}$  is the mean thermal occupation of the mechanical (spin) mode  $\omega_m$  ( $\omega_0$ ) at temperature  $T_k$  of the thermal environments, and  $k_B$  is the Boltzmann constant [54]. The Lindblad superoperators are defined by  $\mathcal{L}[x_k]\rho = (2x_k\rho x_k^\dagger - x_k^\dagger x_k\rho - \rho x_k^\dagger x_k)$ .

#### IV. QUANTUM BATTERY

In this section, we propose to use the spin-mechanical coupling to control the spin so as to store energy in the

harmonic motion of the nanodiamond. The motional modes of a high- $Q$  mechanical oscillator can have very long lifetimes at low pressures and, therefore, are excellent candidates to store coherent energy for extended periods, functioning as a battery.

The spin-mechanical battery protocol involves two key steps, as illustrated in Fig. 2(a):

*Step 1.* A fast resonant  $\pi$  pulse transfers the spin from the ground state  $|0\rangle$  to the excited state  $|-1\rangle$ . After this, the system is allowed to evolve for a time  $\tau_1$ . The pulse switches on the spin-mechanical interaction, which generates a drive on the mechanics of the form  $\mathcal{H}_{\text{int}} \sim g(a + a^\dagger)$ , charging the phonon battery. In a realistic scenario, the spin will decay slightly toward a thermal state due to spin relaxation.

*Step 2.* At time  $\tau_1$ , a 532 nm green light pulse is then used to rapidly reset the spin to the  $|0\rangle$  state, holding it there for a duration  $\tau_2$ , where  $\tau_1 + \tau_2 = 2k\pi/\omega_m$ . This reset operation is nonunitary and, once completed, effectively switches off the mechanical drive, leaving the spin in a product state with the mechanical system.

We focus on the limit where  $\tau_1 \ll \tau_2$ , as decoherence primarily occurs during  $\tau_1$ . To store energy in the battery, we begin with the mechanics in a thermal state and the spin in the  $|0\rangle$  state, i.e.,  $\rho(t=0) = \rho_{\text{th}}(\bar{n}_{\text{th}}) \otimes |0\rangle\langle 0|$ . We then repeatedly apply the battery protocol until the desired amount of energy is stored in the mechanical system. This energy storage is evidenced by a coherent displacement of the thermal state and a corresponding increase in the mean phonon number. To extract the stored energy, one can reverse the overall displacement of the thermal state by applying the control sequence to the spin in a time-reversed manner relative to the original charging sequence. To be effective, the reversed pulse sequence should be applied after a delay of  $2(k+1)\pi/\omega_m$ , following the completion of the charging pulses [Figs. 2(b) and 2(c)].

Beyond indicating energy transfer, the coherent displacement of the mechanical state implies that a portion of the stored energy is available as extractable work. After charging, the mechanical mode occupies a displaced thermal state, which is nonpassive, i.e., the portion of energy extractable via unitary operations. Although our protocol involves nonunitary steps, such as fast optical spin resets, the mechanical state retains coherence that can be harnessed for work extraction by reversing the control sequence or coupling to an external load. Previous works, such as Refs. [67,68], have focused on unitary or thermally isolated dynamics. In contrast, we exploit engineered dissipation as a control resource, enabling repeated, robust charging while maintaining access to extractable energy.

To quantify the extractable portion of the stored energy, we can estimate the coherent energy caused by the displacement  $\mathcal{W}_{\text{coh}}$  of the mechanical oscillator after a single charging pulse. The state of the oscillator is well approximated by a displaced thermal state,  $\rho = D(\alpha)\rho_{\text{th}}D^\dagger(\alpha)$ , whose displacement energy is given by  $\mathcal{W}_{\text{coh}} = \hbar\omega_m|\alpha|^2$ . Ideally, the displacement energy is approximately  $|\alpha| \approx g\tau_1/2$  for small  $\tau$ . Using the parameters from Fig. 2, where  $g/\omega_m = 0.108$  and  $\tau_1 = 2\pi/(128\omega_m)$ , we estimate for a single charging pulse,  $|\alpha| \approx 0.00265$ , corresponding to  $\mathcal{W}_{\text{coh}} \approx 7 \times 10^{-6} \hbar\omega_m$ . While this is a small fraction of the oscillator's energy scale, it nonetheless confirms that even a single charging pulse,

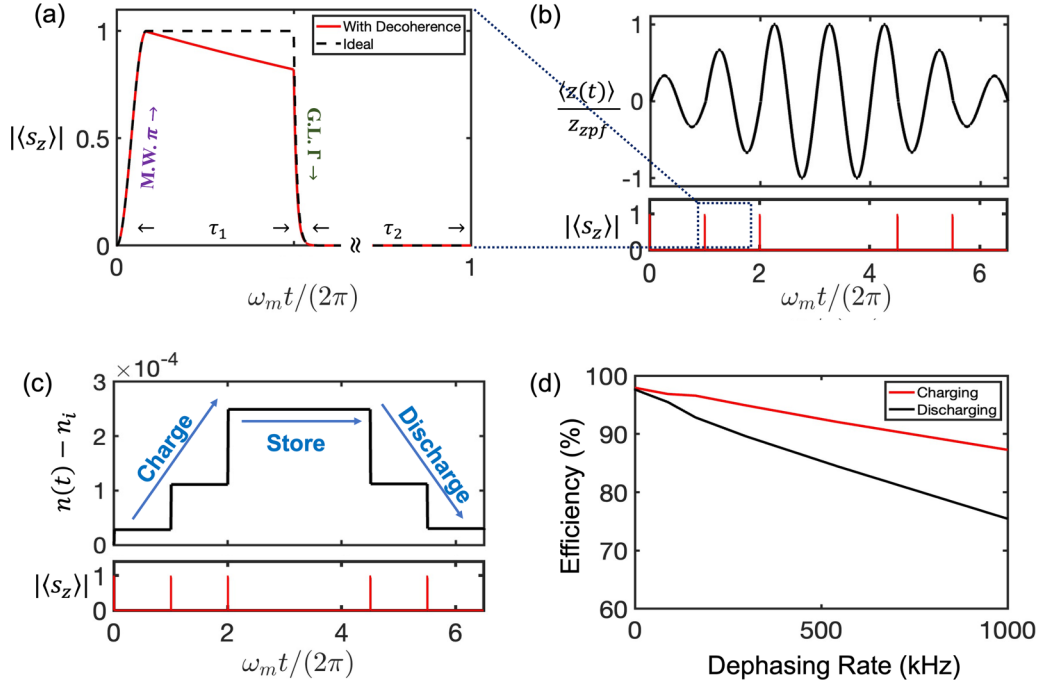


FIG. 2. Schematic of the battery protocol used to store energy in a low-frequency mechanical oscillator via a nitrogen-vacancy (NV) spin. Each cycle begins with the state  $\rho_i \equiv \rho_m^i \otimes |0\rangle\langle 0|$ . A fast resonant microwave (MW  $\pi$ ) pulse transfers the spin to  $|-1\rangle$ , activating the spin-mechanical coupling ( $g/\omega_m = 0.108$ ). After a free-evolution time  $\tau_1$ , a green-laser (GL) pulse resets the spin to  $|0\rangle$ , thereby decoupling it from the oscillator for an additional  $\tau_2$ . This sequence repeats with total period  $\tau_1 + \tau_2 = T_m$ , where  $T_m$  is the oscillator's period. (a)  $\langle s_z \rangle$  vs time; black (red) curves denote spin dynamics without (with) damping at rate  $\gamma_1$ . (b) For  $\tau_1 = T_m/128$ , the spin delivers short “kicks” to the oscillator, incrementally displacing it in phase space. Plotted are the mechanical position  $\langle z \rangle / z_{\text{zpf}}$  (top) and  $\langle s_z \rangle$  (bottom), starting from an initial thermal state with mean occupation  $\bar{n}_i = 2$ . In-phase kicks at  $\tau = 1, 2$  “charge” the oscillator, increasing its amplitude, while no kicks thereafter preserve the stored energy. Out-of-phase kicks at  $\tau = 4.5, 5.5$ , and  $6.5$  “discharge” the oscillator, reducing its amplitude. (c) Evolution of the mean phonon occupation  $n(t) = \langle a^\dagger a \rangle$  showing how each charging, storage, and discharging step, respectively, increases, maintains, or decreases the oscillator's energy  $E_{\text{mech}} = n\hbar\omega_m$ . (d) Charging (red) and discharging (black) efficiencies as functions of the spin dephasing rate (including amplitude damping at  $\gamma_1$ ). We evaluate these efficiencies using Eqs. (8) and (9) in the main text. We observe that the limiting factor for charging and discharging the battery is the spin dephasing rate, which affects discharging efficiency more strongly than charging.

despite dissipation and realistic timing constraints, stores a nonzero amount of coherent, work-extractable energy in the mechanical mode. The amount of work that can be stored unitarily after  $N$  pulses scales as  $N^2 \mathcal{W}_{\text{coh}}$  up to a maximum of  $N \approx Q$ .

We shall now look into the effects of environmental decoherence on the efficiency of the battery. The degree of motional charging is proportional to the duration  $\tau_1$ , and the maximum charging over a single pulse occurs when  $\tau_1 = \pi/\omega_m$ , with the maximum phonon value achieved being  $n_{\text{max}} = (2g/\omega_m)^2$ . For the ideal case without any decoherence, the coherent energy stored closely follows the unitary case,  $\mathcal{W}_{\text{pulse}} \approx \mathcal{W}_{\text{coh}}$ . Adding in the effect of decoherence distorts the pulse shape, as seen in Fig. 2(a); therefore, the energy stored in the mechanical oscillator reduces. We can define the charging efficiency of an uncharged mechanical battery as

$$\eta_{\text{charge}} = \frac{\Delta n_{\text{charge}}}{\Delta n_{\text{max}}}, \quad (8)$$

where  $\Delta n_{\text{charge}}$  is the number of phonons actually gained by the mechanical system due to a single pulse of a duration  $\tau_1$  and  $\Delta n_{\text{max}}$  is the maximum number of phonons gained

due to an ideal pulse (i.e., without dephasing and amplitude damping) of the same duration.

In the case of discharging, we consider a battery that is initially charged by a single pulse. The maximum amount of phonons that can be discharged will be the initially stored phonons  $\Delta n_{\text{charge}}$ . Thereby, we benchmark the discharge efficiency as

$$\eta_{\text{discharge}} = \frac{\Delta n_{\text{discharge}}}{\Delta n_{\text{charge}}}. \quad (9)$$

For the lossless case, the charging and discharging pulse sequences are expected to be symmetrical due to reversibility. In this ideal scenario,  $\Delta n_{\text{charge}}$  should equal  $n_{\text{max}}$ ,  $\Delta n_{\text{discharge}}$  should match  $\Delta n_{\text{charge}}$ , and both the spin and phonon should return to their initial states after a single charge-discharge cycle. However, spin decay, particularly dephasing, impacts the efficiency of these processes. As a result, the system cannot fully return to its original mechanical state, leading to energy loss, as illustrated in Fig. 2(d). Therefore, the maximum amount of retrievable energy stored has an upper limit set by the decoherence processes in the spin system. However, notably, the battery's operation is robust to the initial quantum state of the mechanics, eliminating the need

for a specific initialization of the mechanical oscillator. Unlike previous realizations based on superconducting qubits [69,70], our scheme stores energy in a long-lived mechanical mode via spin-mechanical coupling, offering an alternative route to quantum batteries compatible with room-temperature operation and robust against environmental noise.

## V. COOLING

In this section, we present a scheme to enhance the purity of the low-frequency mechanical mode. Specifically, we demonstrate how to cool the mechanical oscillator, reducing its mean phonon occupation number  $\langle a^\dagger a \rangle$ . This cooling process can be interpreted as an interaction with an effective cold bath at temperature  $T_c$ .

We consider the mechanical system coupled to an NV spin, driven by a continuous microwave field with amplitude  $\Omega$  and detuning  $\Delta$ , as described in Eq. (4). This drive dresses the energy levels of the spin and engineers interactions between the spin and mechanical mode.

To make the interaction apparent, we move to a displaced frame of the mechanics given by  $U(\theta)D(g/\omega_m)$ , where  $D(g/\omega_m) = e^{(g/\omega_m)(a^\dagger - a)}$  and  $U(\theta) = \exp(-i(\theta/2)\sigma_y)$  with  $\theta = \arctan(-2\Omega/\delta)$  and  $\delta = \Delta - 2g^2/\omega_m$ . The Hamiltonian (4) transforms to

$$\bar{\mathcal{H}} = \omega_m a^\dagger a - \frac{\bar{\Delta}}{2} \sigma_z + \frac{g}{2\sqrt{1 + (2\Omega/\delta)^2}} \left( \sigma_z + \frac{2\Omega}{\delta} \sigma_x \right) (a + a^\dagger), \quad (10)$$

where  $\bar{\Delta} = \sqrt{\delta^2 + 4\Omega^2}$  and the dressed basis is expressed as

$$|\downarrow\rangle = -\sin \frac{\theta}{2} |0\rangle + \cos \frac{\theta}{2} |-\rangle, \quad (11)$$

$$|\uparrow\rangle = \cos \frac{\theta}{2} |0\rangle + \sin \frac{\theta}{2} |-\rangle. \quad (12)$$

In this frame, the interaction has two components: a longitudinal coupling,  $\mathcal{H}_\parallel \sim \sigma_z(a + a^\dagger)$ , which results in spin-mechanical Zeeman splitting, and a transverse coupling,  $\mathcal{H}_\perp \sim \sigma_x(a + a^\dagger)$ , which leads to energy exchange between the spin and the mechanical modes. The Jaynes-Cummings Hamiltonian is obtained from Eq. (10) by making the rotating-wave approximation. This approximation becomes valid when  $|\omega_m - \bar{\Delta}| \ll \bar{\Delta} + \omega_m$  and when the coupling is much smaller than the frequencies involved. When we move to the frame rotating with  $H_0 = \omega_m a^\dagger a - \bar{\Delta}/2 \sigma_z$ , we find that there are fast- and slow-rotating terms. In particular, for the parameters we consider below, the longitudinal Hamiltonian rotates ten times faster than the transverse Hamiltonian in this frame. After neglecting the fast-rotating terms and making the rotating-wave approximation (RWA), we obtain

$$\mathcal{H}^{\text{eff}} = \omega_m a^\dagger a - \frac{\bar{\Delta}}{2} \sigma_z + \tilde{g}(a\sigma_+ + a^\dagger\sigma_-), \quad (13)$$

with  $\tilde{g} = (g\Omega/\delta)/(\sqrt{1 + (2\Omega/\delta)^2})$ . We can check the validity of the RWA for the parameters that we will use in the remainder of the paper [given in the caption of Fig. 3(c)], and we find that  $|\omega_m - \bar{\Delta}|/(\bar{\Delta} + \omega_m) \approx 0.706$ , indicating the rotating-wave approximation is marginally valid. We also find  $\tilde{g}/\omega_m \sim 0.073$  and  $\tilde{g}/\bar{\Delta} \sim 0.0125$ , which again confirm the validity of the RWA. In the numerical simulations we perform,

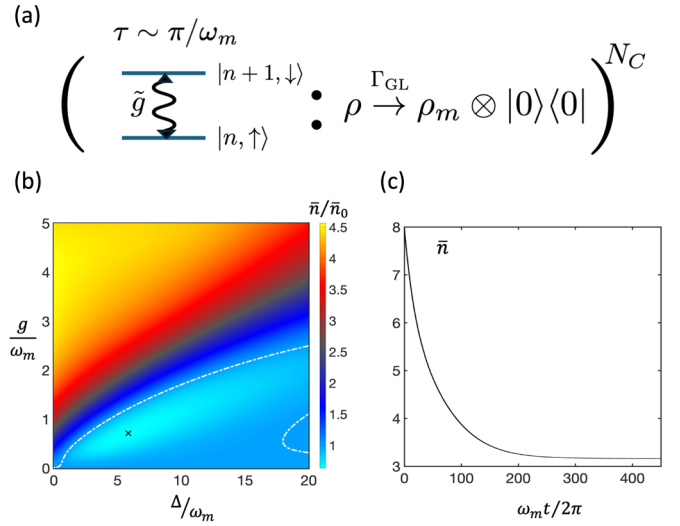


FIG. 3. (a) Illustration of the cooling protocol: The spin subsystem, coupled to a low-frequency mechanical oscillator with strength  $\tilde{g}$ , undergoes repeated operations. Each operation consists of a microwave drive (Rabi frequency  $\Omega$ ) followed by a spin reset using green light at a rate  $\Gamma_{\text{GL}}$ . This sequence efficiently transfers energy from the mechanical mode to the spin, which is then cooled via the reset, reducing the mean phonon occupation number and cooling the mechanical system. The protocol is iterated  $N_c$  times, with each iteration lasting  $\tau \sim \pi/\omega_m$ . (b) Cooling efficiency map: A 2D color map showing the ratio of the mean phonon number  $\bar{n} = \langle a^\dagger a \rangle$  to its initial value  $\bar{n}_0$  after  $N_c = 100$  cooling cycles, plotted as a function of detuning  $\Delta/\omega_m$  and coupling strength  $g/\omega_m$ . The system starts in a thermal state with  $\bar{n}_0 = 8$ . The white dashed line indicates where the ratio equals one. Light blue regions inside this contour highlight optimal cooling areas, where  $\bar{n} < \bar{n}_0$ , demonstrating that precise parameter tuning enhances cooling. (c) Time evolution of phonon occupation: The mean phonon number  $\bar{n}(t)$  is plotted over time for the initial state  $\rho(0) = \rho_{\text{th}}(\bar{n}_0) \otimes |0\rangle\langle 0|$  with  $\bar{n}_0 = 8$ . Parameters used are  $(\Delta, g, \Omega, N_c) = (5.8\omega_m, 0.71\omega_m, \omega_m/2, 900)$ ; other parameters are listed in Table I.

we do not make this approximation and simulate the complete dynamics.

As a result, when the spin is driven in the original frame (4) with a weakly detuned microwave drive, the Hamiltonian in the displaced dressed frame (10) is equivalent to a beam-splitter interaction where energy can be transferred between the phonon and spin.

To efficiently cool the mechanical oscillator, we periodically apply a pulse of green light to reset the spin after a

TABLE I. Parameters used in this paper.

| Quantity                                   | Value                | Unit  |
|--------------------------------------------|----------------------|-------|
| Mechanical frequency, $\omega_m$           | $2\pi \times 50$     | rad/s |
| Quality factor, $Q$                        | 10 000               | —     |
| Dephasing time, $T_2^*$ ( $1/\gamma_2$ )   | $2.5 \times 10^{-3}$ | s     |
| Thermal relaxation, $T_1$ ( $1/\gamma_1$ ) | $5 \times 10^{-3}$   | s     |
| Initialization rate, $\Gamma_{\text{GL}}$  | $10^5$               | Hz    |
| Initialization time, $t_{\text{lase}}$     | $5 \times 10^{-5}$   | s     |

period of beam-splitter interaction. This fast reset process is described by the master equation (5) with the Lindblad superoperator  $\mathcal{L}_{\Gamma_{\text{GL}}}[\sigma_-]$ . We illustrate the cooling process with a diagram in Fig. 3(a).

This beam-splitter interaction serves as the core mechanism through which energy is transferred out of the mechanical mode, achieving cooling. Given that the spin's energy levels are in the MHz–GHz range—significantly higher than the mechanical frequency—we must bring the spin and mechanical oscillator into resonance for efficient energy transfer. This is achieved by driving the spin in the initial frame (4) with a Rabi frequency  $\Omega$ , dressing the spin's energy levels to arrange that the dressed level transition frequency aligns with the mechanical frequency  $\Delta \sim \omega_m$ . The result is an effective beam-splitter Hamiltonian as described by Eq. (13). We maintain the drive for an interaction time  $t_f \sim \pi/\omega_m$ , during which energy is transferred from the mechanical oscillator to the spin, reducing the mean phonon number. Immediately after, the microwave drive is switched off and a green-light pulse is applied to reinitialize the NV center to the  $|0\rangle$  state. This process dissipates the energy gained by the spin into the environment, thereby cooling the spin. By repeating the above steps periodically, we can cool the mechanical mode to the quantum regime incrementally with each cycle. This protocol is repeated  $N_c$  times till the phonon occupation number  $\bar{n} = \langle a^\dagger a \rangle$  reaches a stationary value.

Ideally, resetting the spin to the ground state of the dressed basis will lead to optimal cooling. However, the spin's coherence time  $T_2^*$  poses a limitation, as the states  $|\downarrow\rangle, |\uparrow\rangle$  are tilted in the Bloch sphere, making them susceptible to dephasing. For slow mechanical oscillators, this limitation could hinder cooling efficiency. To mitigate this, we do not prepare the spin to the ground state of the dressed basis after initialization, but we just initialize the spin in the  $|0\rangle$  state and then apply the dressing drive. To determine the optimal parameters for cooling, we numerically sweep the detuning  $\Delta$  and the coupling strength  $g$ . Figure 3(b) reveals a threshold for the coupling strength, beyond which cooling becomes ineffective, irrespective of the detuning. Effective cooling is achieved at low coupling and detuning values, as evidenced by the significant reduction in the phonon number.

## VI. QUANTUM OTTO CYCLE

The Otto cycle is a fundamental thermodynamic cycle used to describe the operation of classical heat engines. It consists of a four-stroke cycle made of adiabatic expansion, isochoric cooling, adiabatic compression, and isochoric heating. By making similar analogies, one can implement an Otto cycle in the quantum regime [3,24]. In this section, we adapt this quantum framework of Otto heat engine to our coupled NV-mechanical oscillator system presented in Fig. 1 by considering the low-frequency phonons as the working fluid.

In traditional heat engines, the heating and cooling parts of the cycle are typically achieved by coupling and decoupling the working fluid with two thermal baths at two temperatures,  $T_h > T_c$ . However, experimentally decoupling the system completely from a heat bath is extremely challenging, which in turn might hamper the efficiency of the engine. Instead, here we consider the working fluid, the phonons of

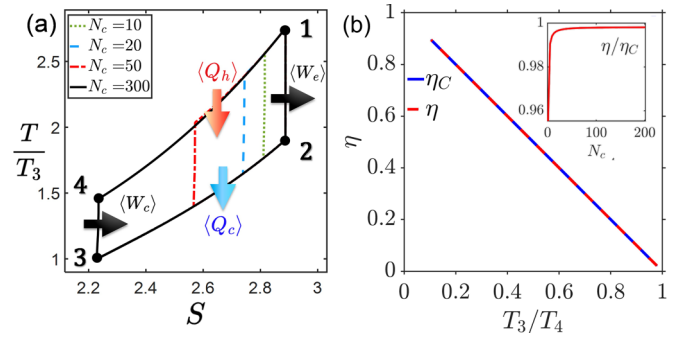


FIG. 4. (a) Temperature-entropy ( $T$ - $S$ ) diagram of the quantum Otto engine for several cooling cycles  $N_c$ . We start with  $\omega_m = 2\pi \times 72$  Hz. Strokes: 1  $\rightarrow$  2—adiabatic decompression  $\omega_m \rightarrow \omega_m(1 - \lambda)$ ; 2  $\rightarrow$  3— $N_c$  spin-mechanical cooling cycles; 3  $\rightarrow$  4—adiabatic compression  $\omega_m(1 - \lambda) \rightarrow \omega_m$ ; 4  $\rightarrow$  1—rethermalization with the motional bath. Arrows indicate the average heats  $\langle Q_h \rangle, \langle Q_c \rangle$  and works  $\langle W_e \rangle, \langle W_c \rangle$ . Unless stated otherwise:  $g = 0.4 \omega_m$ ,  $\omega_m T = 1.45\pi$ ,  $\Delta = 5.8 \omega_m(1 - \lambda)$ , and  $d_t \omega = 1.4 \times 10^4$  rad s $^{-2}$ . (b) Engine efficiency  $\eta$  (solid red) compared with the Otto efficiency [Eq. (15)]  $\eta_c$  (dashed blue) vs  $T_3/T_4$ . All parameters are fixed except the adiabatic stroke time  $\omega_m T \in [0.1\pi, 4.2\pi]$ , resulting in different values of  $\lambda$  (with  $\lambda \equiv \eta_c$ ) and  $T_3/T_4$ . We use initial phonon occupation  $\bar{n}_0 = 5$  and  $N_c = 80$  cooling cycles of duration  $\approx \pi/\omega_m$  each. The straight curve shows that the values of  $\eta$  and  $\eta_c$  are very close. Inset: convergence of  $\eta/\eta_c$  with  $N_c$  (up to 200) at  $\omega_m T = 1.45\pi$ . Additional parameter values are listed in Table I.

the mechanical oscillator, to be in constant thermal contact with the hot bath (at the rate  $\gamma_m$ ) as an “always on” interaction. We will apply the spin-based cooling protocol described in the previous section to engineer an “effective” cold bath during the cooling stroke of the Otto cycle.

Referring to Fig. 4(a), the four strokes of the cycle of our heat engine cycle the mechanical oscillator through the four points in the  $T$ - $S$  diagram, with the Hamiltonian at each point being  $H_1 = \hbar\omega_1 a^\dagger a$ ,  $H_2 = \hbar\omega_2 a^\dagger a$ ,  $H_3 = \hbar\omega_2 a^\dagger a$ , and  $H_4 = \hbar\omega_1 a^\dagger a$ . The strokes of the cycle can then be defined as follows:

**Adiabatic expansion 1  $\rightarrow$  2.** The system starts out in a product state of  $\rho_1 = \rho_{th}(\bar{n}_{th}) \otimes |0\rangle\langle 0|$ , where  $\bar{n}_{th}$  is the mean phonon occupation when in thermal equilibrium with the hot bath. During the adiabatic expansion phase, the mechanical oscillator's frequency is decreased smoothly from its initial value  $\omega_1 = \omega_m$  to a lower value over a time period  $T$ . We consider a linear expansion  $\omega(t) = \omega_m(1 - \lambda \frac{t}{T})$ , where  $\lambda = (d_t \omega) \frac{T}{\omega_m} < 1$  and  $(d_t \omega)$  represents the rate of frequency change. The work done during this stage is  $\langle W_e \rangle = \langle H_2 \rangle - \langle H_1 \rangle < 0$ , where  $\langle H_1 \rangle$  and  $\langle H_2 \rangle$  are the mean energies of the system at the beginning and end of the expansion, respectively.

**Cold isochore 2  $\rightarrow$  3.** In this phase, the system is cooled at the fixed mechanical frequency  $\omega_2 = \omega_m(1 - \lambda)$  by the protocol presented in the previous section to reach an effective temperature  $T_c \equiv T_3$ , which can be preset by choosing values  $(g, \Delta)$  [see Fig. 3(b)]. The heat released by the working fluid during this stage is  $\langle Q_c \rangle = \langle H_3 \rangle - \langle H_2 \rangle$ , where  $\langle H_3 \rangle$  is the mean energy after cooling.

*Adiabatic compression*  $3 \rightarrow 4$ . After the cooling, the mechanical oscillator undergoes compression. The frequency is increased from  $\omega_2$  to  $\omega_1$  over the same time period  $T$ , following the trajectory  $\omega(t) = \omega_m[1 - \lambda(1 - \frac{t}{T})]$ . The work done on the working fluid is  $\langle W_c \rangle = \langle H_4 \rangle - \langle H_2 \rangle > 0$ , where  $\langle H_4 \rangle$  is the energy at the end of the compression.

*Hot isochore*  $4 \rightarrow 1$ . The mechanical oscillator is uncoupled from the spin and freely interacts with the thermal environment, absorbing heat until it reaches, eventually, the initial thermal state, where  $\langle a^\dagger a \rangle = n_{th}$ . The heat exchange is given by  $\langle Q_h \rangle = \langle H_1 \rangle - \langle H_4 \rangle$ . This stage completes the cycle, restoring the system to its initial state and making it ready for the next cycle.

In Fig. 4(a), we plot the temperature-entropy ( $T$ - $S$ ) diagram of the engine for different values of cooling cycle  $N_c$ . The temperature and entropy at each point  $k = 1-4$  along the cycle are computed from  $\beta_k \hbar \omega_k = \log(\frac{\bar{n}_k + 1}{\bar{n}_k})$ , with  $T_k = \frac{1}{k_B \beta_k}$  and  $S(\rho_k) = -\text{Tr}(\rho_k \log \rho_k)$ , where  $\bar{n}_k$  and  $\rho_k$  are the phonon occupation number and density matrix at the  $k$  point, respectively. We note that as  $N_c$  is increased, the cycle can reach lower  $S$  values as the temperature of the mechanical oscillator decreases, slowly coming into equilibrium with the temperature of the effective cold bath,  $T_c$ . The form of the cycle is similar to the classical Otto cycle from which four temperatures are defined.

To quantize the efficiency of the heat engine, we defined, similarly to classical heat engines, the quantum efficiency  $\eta$  of the cycle. It is defined as the ratio of net work output to the heat absorbed from the hot bath:

$$\eta = 1 + \frac{\langle Q_c \rangle}{\langle Q_h \rangle} = -\frac{(\langle W_e \rangle + \langle W_c \rangle)}{\langle Q_h \rangle}. \quad (14)$$

For an ideal Otto cycle, the theoretical quantum efficiency limit can be expressed in terms of the frequencies as

$$\eta_c = 1 - \frac{\omega_2}{\omega_1} = \lambda = 1 - \frac{T_3}{T_4}. \quad (15)$$

In Fig. 4(b), we plot the values of  $\eta$  and  $\eta_c$  as a function of the ratio  $T_3/T_4$  for quite a high value of  $N_c$  and observe that the computed efficiency is very close to the ideal Otto efficiency. In the inset, we explore this more by varying  $N_c$  and observe that the computed efficiency is a very rapid function of  $N_c$ .

This spin-mechanical quantum heat engine demonstrates an efficiency approaching the ideal Otto efficiency without the need for quasistatic operation.

While the mechanical mode remains in contact with the hot bath, the key to achieving adiabaticlike behavior lies in the distinct timescales involved. In our proposal, owing to the high  $Q$  of the resonator, the coupling between the thermal bath and the system is very weak, resulting in a long thermalization time. The time durations for the compression and expansion strokes are chosen to be short enough to prevent thermal equilibrium with the bath. At the same time, the timescales for these expansion and compression strokes are carefully

selected to avoid being so rapid that they excite diabatic transitions.

## VII. DISCUSSION AND CONCLUSION

We have proposed two quantum thermal machines—a quantum battery and a quantum heat engine—built upon a single nitrogen-vacancy (NV) center embedded in a harmonically trapped nanodiamond. By exploiting the NV spin's controllability and its coupling to its low frequency yet long-lived mechanical mode, we demonstrated that coherent energy exchange can be achieved even when the mechanical frequency is far below the spin's transition and relaxation rates.

These protocols can be implemented across a range of spin-mechanical platforms, including cantilevers and levitated oscillators [61,64,66]. For a magnetic-field gradient  $G_z \approx 10^3$  T/m, mechanical oscillator frequency  $\omega_m \approx 2\pi \times 50$  Hz and a radius of diamond  $r \approx 1$   $\mu$ m corresponding to a mass of  $m \approx 10^{-11}$  kg, we can achieve a coupling  $g/\omega_m$  of the order 0.05–0.1, the regime where we can optimally cool the mechanical oscillator as shown in Ref. [64,66]. In this regime, the mechanical oscillator can operate both as a “quantum battery”—storing or releasing energy on demand—and as the working fluid of a heat engine.

While NV spins typically exhibit microsecond-scale dephasing times  $T_2^*$ , dynamical decoupling methods can extend coherence to milliseconds or more [71,72]. Alternatively, we can achieve a higher coherence time by leveraging the nuclear spins in diamonds. Recent demonstrations have shown control over nuclear spin degrees of freedom in levitated microdiamonds [73].

Our approach does not require specialized preparation of the mechanical oscillator; it readily accommodates thermal mechanical states of oscillators of a wide range of frequencies. Consequently, the results presented here highlight a practical pathway toward quantum thermodynamic devices capable of high coherence, efficient energy storage, and controlled heat exchange.

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## DATA AVAILABILITY

All the codes to reproduce the results in the research can be found here in Refs. [74,75].

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- [74] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/mgjh-rws9> for python and MATLAB codes to simulate and reproduce the Figs. 2–4.
- [75] <https://github.com/Mohatifi/Spin-mechanical-thermal-machines>.