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Nonreciprocal Quantum Batteries

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Nonreciprocity, arising from the breaking of time-reversal symmetry, has become a fundamental tool in diverse quantum technology applications. It enables directional flow of signals and efficient noise suppression, constituting a key element in the architecture of current quantum information and computing systems. Here we explore its potential in optimizing the charging dynamics of a quantum battery. By introducing nonreciprocity through reservoir engineering during the charging process, we induce a directed energy flow from the quantum charger to the battery, resulting in a substantial increase in energy accumulation. Despite local dissipation, the nonreciprocal approach demonstrates a fourfold increase in battery energy compared to conventional charger-battery systems. This effect is observed in the stationary limit and remains applicable even in overdamped coupling regimes, eliminating the need for precise temporal control over evolution parameters. Our result can be extended to a chiral network of quantum nodes, serving as a multicell quantum battery system to enhance storage capacity. The proposed approach is straightforward to implement using current state-of-the-art quantum circuits, both in photonics and superconducting quantum systems. In a broader context, the concept of nonreciprocal charging has significant implications for sensing, energy capture, and storage technologies or studying quantum thermodynamics.

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Nonreciprocal devices play an important role in optimizing system performance by suppressing undesired signal paths, noises, and spurious modes [1]. The discreet foundation of nonreciprocity lies in breaking time-reversal symmetry, a fundamental principle governing electromagnetic wave behavior [2,3]. The intentional breaking of time-reversal symmetry emerges as a strategy in engineering nonreciprocal devices, providing precise control over signal pathways and improving the efficiency of quantum devices. Nonreciprocity not only facilitates selective signal transmission in one direction while hindering the opposite to prevent interference but also stimulates exploration into inventive quantum device functionalities. In circuit quantum electrodynamics [4,5] and nanophotonics [6–9] quantum communications, circulators serve as essential components, operating as single-port couplers or isolators. Their critical role lies in shielding vulnerable quantum states of cavities and qubits against electromagnetic noises and back reflection of intense signals or pumps.

In recent years, there has been a growing interest in both theoretical exploration and experimental application of nonreciprocity in a wide range of systems, spanning from microwave quantum circuits [10–19] to photonic devices [9,20,21] and opto-electromechanical systems [22–29]. This interest is directed towards the design of efficient optical and microwave isolators and the investigation of various manifestations of nonreciprocity. However, the importance of nonreciprocity goes beyond merely suppressing unwanted signals. For instance, nonreciprocal devices can be utilized to regulate the flow of thermal noise, enabling the realization of a thermal rectifier in nanoscale quantum devices [30]. In this Letter, we study a novel aspect of nonreciprocal interactions: the role of nonreciprocity in facilitating efficient energy transfer and storage between quantum systems. We show that this concept has the potential for the implementation of nonreciprocal quantum batteries.

The exploration of utilizing quantum systems as quantum batteries for capturing and storing energy has recently gained significant attention [31–45]. The primary aim of quantum batteries is to enhance the efficiency of the energy storage and charging power [32,33,36,39,41,43], utilizing quantum resources such as entanglement [39,46,47], quantum optimal control [48,49], and quantum catalysis processes [50]. Using interconnected quantum nodes or units offers the potential to design multicell quantum batteries, enhancing the capacity of energy storage. However, attempting to establish such a system solely through coherent interaction between cells is impractical. As with any other quantum systems, the challenges of decoherence and losses present a potential threat to preserving the essential quantum properties necessary for constructing reliable quantum batteries, as demonstrated experimentally [40,51–53].

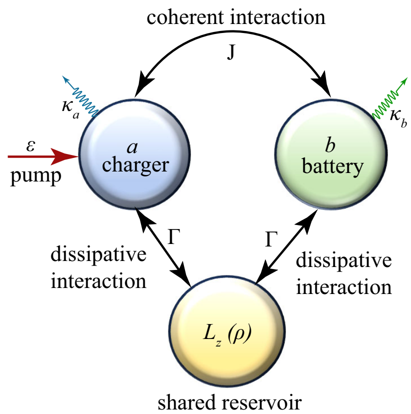

FIG. 1. Schematic representation of the nonreciprocal quantum battery system. A quantum charger a coherently interacts with a quantum battery b with a coupling rate J. The system's energy is supplied by an external pump with frequency ω_L and amplitude $\mathcal{E}$. Both the charger and the battery simultaneously couple to a shared reservoir with a rate Γ. In this configuration, energy directly flows from the charger to the battery, while the nonreciprocal condition effectively suppresses energy backflow. Here, κ_a and κ_b describe the local damping rates of each mode.

We introduce a novel approach designed to maximize energy storage in a quantum battery while minimizing energy dissipation to the surrounding environment during the charging process. Our method utilizes reservoir engineering [54–56], where a carefully designed dissipative environment facilitates efficient energy transfer between the modes of interest. Incorporating a dissipative reservoir, like an auxiliary waveguide, initiates an effective dissipative interaction between the charger and the battery. To create a nonreciprocal energy flow, we utilize an interference-like process, carefully balancing the induced dissipative interaction with its coherent counterpart. This nonreciprocal energy transfer boosts the accumulation of energy in the quantum battery, enhancing energy storage capabilities and reducing energy backflow to the charger. Our nonreciprocal quantum battery can be implemented in opto-electromechanical systems where a microwave resonator can act as a charger and a mechanical resonator as a battery [28,29]. Additionally, alternative platforms such as nonreciprocal superconducting [57,58] or magnonic [59,60] circuits can be used for the experimental realization of our proposed system. Our results have the potential to contribute to the advancement of understanding and engineering nonreciprocal energy transfer in quantum systems. Specifically, this opens up a new avenue for enhancing energy storage efficiency in micro- and nanoscaled quantum devices for future applications.

Figure 1 illustrates the system's schematic, where a harmonic oscillator serves as the charger with a resonance frequency of ω_a and local damping rate κ_a. This charger interacts with another mode acting as the battery, which has a frequency of ω_b and a damping rate κ_b. The interaction between the charger and the battery is established through a coherent coupling with a rate of J. A classical drive field with a frequency of ω_L and amplitude $\mathcal{E}$ is utilized to provide energy to the charger. This energy then will be used to charge the battery. The Hamiltonian describing this driven bipartite system can be written as ($\hbar = 1$)

$$H = \omega_a a^\dagger a + \omega_b b^\dagger b + (J a^\dagger b + J^* b^\dagger a) + \mathcal{E}(e^{i\omega_L t} a + e^{-i\omega_L t} a^\dagger), \tag{1}$$

where a and b are the annihilation operators of the charger and the battery, respectively. For simplicity, we assume that both the charger and the battery have the same resonance frequency, $\omega = \omega_a = \omega_b$, see Supplemental Material [61].

Until now, the model remains similar to the charger-battery model previously studied in other works [37,39,62], where the reciprocal interaction between the charger and battery results in the exchange of energy between them. To achieve nonreciprocity, we additionally consider the existence of a dissipative interaction between the charger and the battery modes. This interaction occurs when both modes are coupled to a common (shared) nonlocal reservoir, characterized by a coupling rate Γ. By adiabatically eliminating the reservoir, an effective dissipative coupling between the charger and battery is established [56]. In our analysis, we assume that the reservoir is Markovian and thus, we employ the standard master equation to describe the dynamics of the system $\dot{\rho} = -i[H, \rho] + \sum_{i=a,b} \kappa_i \mathcal{L}_i[\rho] + \Gamma \mathcal{L}_z[\rho]$, where $\mathcal{L}_c[\rho] = c\rho c^\dagger - \frac{1}{2}\{c^\dagger c, \rho\}$ represents the dissipative superoperator resulting from the coupling to the shared reservoir. Here, $z = p_a a + p_b b$ with p_a and p_b describe the coupling of the charger and battery, respectively, to the shared reservoir [56]. For simplicity, we are considering a Markovian reservoir, but the nonreciprocity can be implemented in non-Markovian regimes as well [63]. The first term of the master equation describes the coherent coupling between the charger and battery, while the second term describes the local damping of each mode to their bath with rate $\kappa_{a/b}$. The last term represents the dissipation of both modes into the shared reservoir, resulting in dissipative coupling between the two modes. This coupling can be realized by treating the shared bath as a damped cavity, waveguide, or transmission line [29,59,60].

In the Schrödinger picture, the evolution of the operators is governed by the following equations:

$$\frac{d\langle a\rangle}{dt} = -\left(\frac{\Lambda_a}{2} + i\omega\right)\langle a\rangle - i\left(J + i\mu\frac{\Gamma}{2}\right)\langle b\rangle - ie^{-i\omega_L t}\mathcal{E},$$
$$\frac{d\langle b\rangle}{dt} = -\left(\frac{\Lambda_b}{2} + i\omega\right)\langle b\rangle - i\left(J^* + i\mu^*\frac{\Gamma}{2}\right)\langle a\rangle. \tag{2}$$

However, the energy stored in the battery is given by $E_B = \mathrm{Tr}_a[\rho H] = \omega_b \langle b^\dagger b\rangle$ where $\mathrm{Tr}_a[..]$ is the partial trace over the charger's degree of freedom. Thus, calculating the energy of the system requires determining the second moments of the operators

$$\frac{d\langle a^\dagger a\rangle}{dt} = -\Lambda_a\langle a^\dagger a\rangle - 2\mathrm{Re}\left\{i\left(J + i\mu\frac{\Gamma}{2}\right)\langle a^\dagger b\rangle\right\} - 2\mathrm{Im}\left\{e^{i\omega_L t}\mathcal{E}\langle a\rangle\right\},$$
$$\frac{d\langle b^\dagger b\rangle}{dt} = -\Lambda_b\langle b^\dagger b\rangle + 2\mathrm{Re}\left\{i\left(J - i\mu\frac{\Gamma}{2}\right)\langle a^\dagger b\rangle\right\}, \quad (3)$$

which explains that the population of each mode is directly influenced by $\langle a^\dagger b\rangle$, indicating the excitation hopping between the two modes, expressed by

$$\frac{d\langle a^\dagger b\rangle}{dt} = -\left(\frac{\Lambda_a + \Lambda_b}{2}\right)\langle a^\dagger b\rangle - i\left(J^* + i\mu^*\frac{\Gamma}{2}\right)\langle a^\dagger a\rangle + i\left(J^* - i\mu^*\frac{\Gamma}{2}\right)\langle b^\dagger b\rangle + ie^{i\omega_L t}\mathcal{E}\langle b\rangle, \quad (4)$$

where $\mu = -p_b p_a^*$, $\Gamma_i = \Gamma|p_i|^2$, and $\Lambda_i = \Gamma_i + \kappa_i$ with $i = a, b$. As long as both p_a and p_b are nonzero, we can always rescale them such that $|\mu| = 1$, as the remaining factor is absorbed into Γ.

The nonreciprocity now can be achieved by balancing the dissipative and coherent interaction rates. In Eqs. (2)–(4), by selecting specific values for the coupling $J = -i\mu(\Gamma/2)$, the behavior of $\langle a\rangle$ and $\langle a^\dagger a\rangle$ remains entirely unaffected by the presence of the battery which implies a unidirectional flow of energy from the charger to the battery, without any energy backflow. This imposes nonreciprocal conditions and ensures an effective accumulation of energy in the battery. Solving the complete set of Eqs. (2)–(4), we obtain the stored energy in the battery in the nonreciprocal regime:

$$E_B^{\mathrm{nr}}(t) = \frac{16\omega\Gamma^2\mathcal{E}^2}{\Lambda_a^2\Lambda_b^2}\left(1 + \left[\frac{\Lambda_a e^{-\frac{1}{2}\Lambda_b t} - \Lambda_b e^{-\frac{1}{2}\Lambda_a t}}{(\Lambda_a - \Lambda_b)^2}\right] \times \left[(\Lambda_a e^{-\frac{1}{2}\Lambda_b t} - \Lambda_b e^{-\frac{1}{2}\Lambda_a t}) - 2(\Lambda_a - \Lambda_b)\right]\right), \quad (5)$$

where we considered the drive to be in resonance with each local mode $\delta = \omega_L - \omega = 0$, see Supplemental Material [61] for the derivation and proof of optimality of the resonant drive. In the limit of symmetric damping rates, i.e., $\Lambda = \Lambda_a = \Lambda_b$ (which requires $\Gamma = \Gamma_a = \Gamma_b$ and $\kappa_a = \kappa_b$), Eq. (5) can be simplified to

$$E_B^{\mathrm{nr}}(t) = \frac{16\omega\Gamma^2\mathcal{E}^2}{\Lambda^4}\left[1 + e^{-\Lambda t/2}(e^{-\Lambda t/2} - 2)\right], \quad (6)$$

where for $t \to \infty$, this function reaches its steady state value $E_B^{\mathrm{nr}}(\infty) = (16\omega\Gamma^2\mathcal{E}^2/\Lambda^4)$.

We can also calculate the energy of the charger under nonreciprocal condition

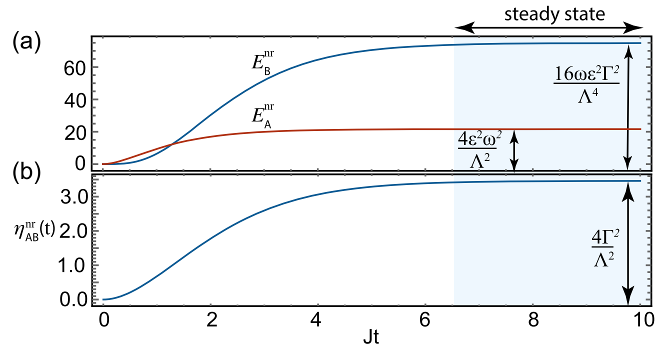

FIG. 2. (a) The total energy of the charger E_A^{nr} and battery E_B^{nr} plotted against the scaled time Jt in the nonreciprocal regime $|J| = |\Gamma/2|$. This plot illustrates that as the system approaches a steady state, the energy of the battery surpasses that of the charger. (b) The ratio between the energies of the battery and charger, $\eta_{AB}^{\mathrm{nr}}(t) = E_B^{\mathrm{nr}}/E_A^{\mathrm{nr}}$. In the steady-state limit, this ratio converges to $\mathcal{C}_d = 4\Gamma^2/\Lambda^2$. Here, we consider $\Gamma = \Gamma_a = \Gamma_b = 0.04\omega$, $\mathcal{E} = 0.1\omega$, $\kappa_a = \kappa_b = 0.003\omega$, and $J = \Gamma/2$.

$$E_A^{\mathrm{nr}}(t) \coloneqq \omega\langle a^\dagger a\rangle = \frac{4\omega\mathcal{E}^2\left(1 - e^{-\frac{1}{2}\Lambda t}\right)^2}{\Lambda^2}, \quad (7)$$

with steady state value $E_A^{\mathrm{nr}}(\infty) = (4\omega\mathcal{E}^2/\Lambda^2)$. In Fig. 2(a), the energies of the battery E_B^{nr} and charger E_A^{nr} are plotted against the scaled time Jt in the nonreciprocal regime. Initially, the energy stored in the battery is lower than that in the charger, but it progressively increases and reaches its steady-state value. The ratio of the energies of the charger and battery $\eta_{AB}^{\mathrm{nr}}(t) = [E_B^{\mathrm{nr}}(t)/E_A^{\mathrm{nr}}(t)]$ provides additional insights into the energy distribution between the two modes. In the steady-state limit, we find that $\eta_{AB}^{nr}(\infty) = \mathcal{C}_d$, where $\mathcal{C}_d = 4\Gamma^2/\Lambda^2$ represents dissipative cooperativity between the charger and the battery through the shared reservoir. Efficient energy storage in the battery in the nonreciprocal scenario requires $\mathcal{C}_d > 1$ or, equivalently, $\Gamma > \kappa$, where $\kappa = \kappa_a = \kappa_b$. Therefore, as long as the coupling of both the charger and the battery to the shared reservoir surpasses the local damping of each mode, the energy of the battery exceeds that of the charger. Figure 2(b) illustrates the parameter $\eta_{AB}^{\mathrm{nr}}(t)$ versus the scaled time Jt. This ratio surpasses 1 immediately after the interaction is turned on and reaches its steady-state value set by $\mathcal{C}_d$.

The efficiency of the nonreciprocal charging method can be examined by comparing the battery's energy to that of a conventional reciprocal charger-battery system [35,37]. This comparison can be done by removing the shared reservoir and setting $\Gamma_i = \Gamma = 0$ in Eqs. (2)–(4), resulting in

$$E_B = \zeta\left(1 - \alpha(t)e^{-\kappa_{ab}t/4} + \left[\beta(t) - \frac{2\kappa_a\kappa_b(\mathcal{C}+1)}{\Delta^2}\right]e^{-\kappa_{ab}t/2}\right), \quad (8)$$

with $\zeta = [4\omega\mathcal{E}^2\mathcal{C}/\kappa_a\kappa_b(\mathcal{C}+1)^2]$ and

$$\alpha(t) = \frac{2}{\Delta}(\Delta \cosh \Delta t/4 + \kappa_{ab} \sinh \Delta t/4),$$
$$\beta(t) = \frac{(\Delta^2 + \kappa_{ab}^2)\cosh \Delta t/2 - 2\kappa_{ab}\Delta \sinh \Delta t/2}{2\Delta^2}, \quad (9)$$

where $\Delta = \sqrt{-16|J|^2 + (\kappa_a - \kappa_b)^2}$, $\kappa_{ab} = \kappa_a + \kappa_b$, and $\mathcal{C} = 4J^2/\kappa_a\kappa_b$ is the coherent interaction cooperativity between charger and battery.

In the underdamped regime when $4|J| > \{\kappa_a, \kappa_b\}$ or for $4|J| > \{|\kappa_a - \kappa_b|\}$, Δ is imaginary. In this case, the hyperbolic functions in Eqs. (9) transform into oscillating terms, representing the energy exchange between the battery and charger. These oscillations experience exponential decay due to the local damping in each mode with rate $\kappa_{a/b}$, leading to a slow and inefficient charging process. When the excitation returns to the charger, a portion of it leaks into the charger's environment before reaching back to the battery. In the nonreciprocal scenario, the dynamics of the charger is entirely separate from the battery. This isolation eliminates time-dependent oscillations in the battery's energy, enabling a more efficient and faster charging process. In the following section, we first explore the underdamped regime (Δ is imaginary in general), comparing energy storage in the battery under nonreciprocal and reciprocal conditions. Subsequently, we focus on real values of Δ in the overdamped regime. In both regimes, we identify an optimal condition wherein nonreciprocity consistently leads to efficient energy storage in the battery.

In Fig. 3(a), the comparison of the battery's energy with and without nonreciprocity is presented over the scaled time Jt for $\kappa_a = \kappa_b$, where Δ is imaginary (underdamped, i.e., $4|J| > \{\kappa_a, \kappa_b\}$). At a finite time limit, the first peak in the energy of the battery occurs more rapidly in the reciprocal case, resulting in a higher charging power compared to the nonreciprocal conditions, as shown in Fig. 3. However, the system reaches to its steady state much faster in the nonreciprocal regime with much higher stored energy compared to the reciprocal charger-battery system. In general, however, the coupling of the charger and battery to the shared reservoir allows for the optimization of system parameters, as explained in the Supplemental Material [61].

The impact of nonreciprocity and its advantages becomes more apparent when examining the parameter $\eta_{BB}(t) = [E_B^{\rm nr}(t)/E_B(t)]$, as seen in Fig. 3(b). This function oscillates and eventually reaches its steady state, set by

$$\eta_{BB}(\infty) = 4\left(\frac{1+\mathcal{C}}{(\sqrt{\mathcal{C}}+\xi)(\sqrt{\mathcal{C}}+\frac{1}{\xi})}\right)^2, \quad (10)$$

where $\xi = \sqrt{\kappa_a/\kappa_b}$ and for a meaningful comparison, we use the nonreciprocity condition and set $|\Gamma| = 2|J|$. When $\mathcal{C} > \xi^2$ or $\mathcal{C} < 1/\xi^2$, the nonreciprocal regime yields higher stored energy in the battery compared to the reciprocal charger-battery regime. When $\kappa_a = \kappa_b$ or $\xi = 1$, we observe $\eta \geq 1$. In such cases, for $\mathcal{C} \ll 1$ or $\mathcal{C} \gg 1$, the energy stored in the battery during the nonreciprocal regime is four times greater than that of a conventional charger-battery system $\eta_{BB}(\infty) \approx 4$.

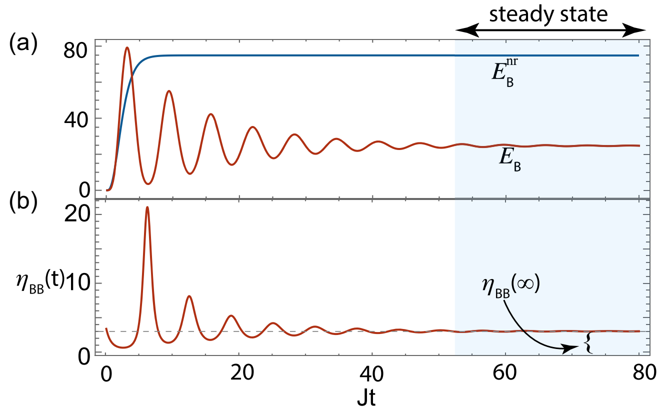

FIG. 3. (a) The energy of the battery in reciprocal $E_B(t)$ and nonreciprocal $E_B^{\rm nr}(t)$ regimes plotted versus the scaled time Jt for imaginary Δ when $4|J| > \{\kappa_a, \kappa_b\}$ or $4|J| > \{|\kappa_a - \kappa_b|\}$. (b) The parameter $\eta_{BB}(t) = E_B^{\rm nr}(t)/E_B(t)$ versus the scaled time Jt. In the steady-state limit, this ratio converges to $\eta_{BB}(\infty)$. In both figures, we consider $\Gamma = \Gamma_a = \Gamma_b = 0.04\omega$, $\mathcal{E} = 0.1\omega$, $\kappa_a = \kappa_b = 0.003\omega$, and $J = \Gamma/2$, leading to $\Delta/\omega = 0.08i$.

The result, however, can be extended to the overdamped regime or real values of Δ where $4|J| < \{\kappa_a, \kappa_b, |\kappa_a - \kappa_b|\}$. Figure 4(a), compares $E_B(t)$ and $E_B^{\rm nr}(t)$ versus scaled time Jt for $\kappa_a > \kappa_b$ and $4J^2 < |\kappa_a - \kappa_b|^2$. As seen, the energy stored in the battery achieves its maximum without oscillations in both reciprocal and nonreciprocal cases. In this scenario, the reciprocal charger-battery approach yields superior results compared to the nonreciprocal charging mechanism. However, by optimizing the system parameters and setting $\Gamma_a = \xi\Gamma$ and $\Gamma_b = \Gamma/\xi$, nonreciprocity can result in more efficient energy storage $E_{B,\rm opt}^{\rm nr}(t)$ in the battery, see Fig. 4(a).

In Fig. 4(b), we plot the $\eta_{BB}(t)$ with and without optimization, confirming that the optimization leads to increased energy accumulation in the battery. In the steady state, we obtain

$$\eta_{BB}^{\rm opt}(\infty) = \frac{\Gamma^2(4J^2 + \kappa_a\kappa_b)^2}{J^2(\Gamma + \sqrt{\kappa_a\kappa_b})^4}, \quad (11)$$

and by setting $\Gamma = 2J$ we get

$$\eta_{BB}^{\rm opt}(\infty) = 4\left(\frac{1+\mathcal{C}}{(\sqrt{\mathcal{C}}+1)^2}\right)^2, \quad (12)$$

in which for $\mathcal{C} \gg 1$ or $\mathcal{C} \ll 1$, $\eta_{BB}^{\rm opt} \approx 4$, see Fig. 4(c). We note that this result is general, and as long as the optimal condition holds ($\Gamma_a = \xi\Gamma$ and $\Gamma_b = \Gamma/\xi$), nonreciprocity

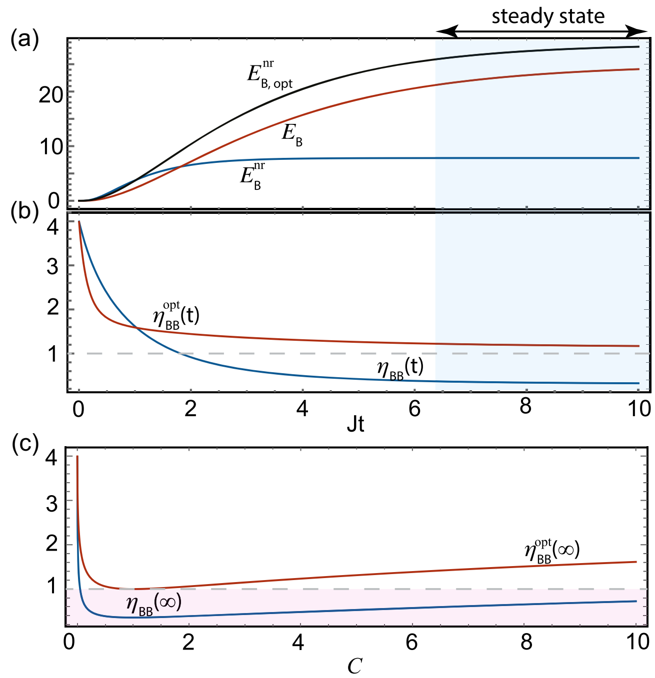

FIG. 4. Comparing the energy of the battery in reciprocal $E_B(t)$, nonreciprocal $E_B^{\rm nr}(t)$, and optimized nonreciprocal $E_{B,\rm opt}^{\rm nr}(t)$ regimes, presented against the scaled time Jt. (b) A comparison of the parameter $\eta_{BB}(t)$ and its optimized counterpart $\eta_{BB}^{\rm opt}(t)$ over time. (c) The steady-state values of parameters $\eta_{BB}(\infty)$ and $\eta_{BB}^{\rm opt}(\infty)$ versus cooperativity $\mathcal{C}$. For (a) and (b), we assume $\Gamma = \Gamma_a = \Gamma_b = 0.01\omega$, $\mathcal{E} = 0.1\omega$, $\kappa_a = 0.1\omega$, $\kappa_b = 0.003\omega$, and $J = \Gamma/2$, corresponding to $\Delta \approx 0.095\omega$. In (c), optimization conditions are applied with $\Gamma_a = \xi\Gamma$ and $\Gamma_b = \Gamma/\xi$, where $\xi = \sqrt{\kappa_a/\kappa_b}$.

consistently outperforms the reciprocal charger-battery system.

In this Letter, we theoretically proposed and investigated a nonreciprocal approach to enhance energy storage in a quantum battery. We demonstrated how breaking reciprocity via reservoir engineering enables unidirectional energy flow, reducing energy dissipation into the surroundings, particularly advantageous for higher dissipation rates. The underlying physics relies on an interference-like phenomenon where coherent coupling between the charger and battery counteracts dissipative interaction. Here we did not discuss the work extracted from the battery or study ergotropy. The reason for that is when a quantum harmonic oscillator is linearly coupled to a bath at zero temperature, Markovian dynamics map a coherent state into another coherent state [64], implying that all internal energy of the battery can be extracted as ergotropy.

Our proposed approach can be implemented using current state-of-the-art quantum photonic systems or microwave superconducting circuits. Moreover, this system is scalable to a chain of batteries, forming a chiral network of quantum cells. The potential benefits include enhanced energy storage efficiency and capacity, offering a pathway for realistic quantum batteries in nano- or microscaled quantum devices. In the broader context of advancing quantum technology with more complex systems and numerous qubits, quantum batteries stand out as potentially integral components. Their usage can extend to enabling reversible logic gates in quantum computing processors [65–67]. Beyond that, our research introduces novel possibilities, such as leveraging nonreciprocity for the study of quantum thermodynamic phenomena and understanding energy flow in quantum devices. Additionally, it could spark the exploration of quantum batteries in naturally nonreciprocal materials or magneto-optic systems.

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[1] D. Jalas, A. Petrov, M. Eich, W. Freude, S. Fan, Z. Yu, R. Baets, M. Popovic, A. Melloni, J. D. Joannopoulos, M. Vanwolleghem, C. R. Doerr, and H. Renner, What is—and what is not—an optical isolator, Nat. Photonics **7**, 579 (2013).
[2] P. Lodahl, S. Mahmoodian, S. Stobbe, A. Rauschenbeutel, P. Schneeweiss, J. Volz, H. Pichler, and P. Zoller, Chiral quantum optics, Nature (London) **541**, 473 (2017).
[3] C. Caloz, A. Alù, S. Tretyakov, D. Sounas, K. Achouri, and Z.-L. Deck-Léger, Electromagnetic nonreciprocity, Phys. Rev. Appl. **10**, 047001 (2018).
[4] B. A. Auld, The synthesis of symmetrical waveguide circulators, IRE Trans. Microw. Theory Techn. **7**, 238 (1959).
[5] D. M. Pozar, *Microwave Engineering*, 4th ed. (John Wiley & Sons, Inc., 2012).
[6] E. Li, B. J. Eggleton, K. Fang, and S. Fan, Photonic Aharonov-Bohm effect in photon-phonon interactions, Nat. Commun. **5**, 3225 (2014).
[7] H. Lira, Z. Yu, S. Fan, and M. Lipson, Electrically driven nonreciprocity induced by interband photonic transition on a silicon chip, Phys. Rev. Lett. **109**, 033901 (2012).
[8] K. Fang, Z. Yu, and S. Fan, Photonic Aharonov-Bohm effect based on dynamic modulation, Phys. Rev. Lett. **108**, 153901 (2012).
[9] S. Manipatruni, J. T. Robinson, and M. Lipson, Optical nonreciprocity in optomechanical structures, Phys. Rev. Lett. **102**, 213903 (2009).

[10] A. Kamal, J. Clarke, and M. H. Devoret, Noiseless non-reciprocity in a parametric active device, Nat. Phys. **7**, 311 (2011).

[11] B. Abdo, K. Sliwa, L. Frunzio, and M. Devoret, Directional amplification with a Josephson circuit, Phys. Rev. X **3**, 031001 (2013).

[12] N. A. Estep, D. L. Sounas, J. Soric, and A. Alu, Magnetic-free non-reciprocity and isolation based on parametrically modulated coupled-resonator loops, Nat. Phys. **10**, 923 (2014).

[13] G. Viola and D. P. DiVincenzo, Hall effect gyrators and circulators, Phys. Rev. X **4**, 021019 (2014).

[14] K. M. Sliwa, M. Hatridge, A. Narla, S. Shankar, L. Frunzio, R. J. Schoelkopf, and M. H. Devoret, Reconfigurable Josephson circulator/directional amplifier, Phys. Rev. X **5**, 041020 (2015).

[15] J. Kerckhoff, K. Lalumière, B. J. Chapman, A. Blais, and K. W. Lehnert, On-chip superconducting microwave circulator from synthetic rotation, Phys. Rev. Appl. **4**, 034002 (2015).

[16] F. Lecocq, L. Ranzani, G. A. Peterson, K. Cicak, R. W. Simmonds, J. D. Teufel, and J. Aumentado, Nonreciprocal microwave signal processing with a field-programmable Josephson amplifier, Phys. Rev. Appl. **7**, 024028 (2017).

[17] A. C. Mahoney, J. I. Colless, S. J. Pauka, J. M. Hornibrook, J. D. Watson, G. C. Gardner, M. J. Manfra, A. C. Doherty, and D. J. Reilly, On-chip microwave quantum Hall circulator, Phys. Rev. X **7**, 011007 (2017).

[18] E. I. Rosenthal, B. J. Chapman, A. P. Higginbotham, J. Kerckhoff, and K. W. Lehnert, Breaking Lorentz reciprocity with frequency conversion and delay, Phys. Rev. Lett. **119**, 147703 (2017).

[19] B. J. Chapman, E. I. Rosenthal, J. Kerckhoff, B. A. Moores, L. R. Vale, J. A. B. Mates, G. C. Hilton, K. Lalumière, A. Blais, and K. W. Lehnert, Widely tunable on-chip microwave circulator for superconducting quantum circuits, Phys. Rev. X **7**, 041043 (2017).

[20] L. Bi, J. Hu, P. Jiang, D. H. Kim, G. F. Dionne, L. C. Kimerling, and A. RossC., On-chip optical isolation in monolithically integrated non-reciprocal optical resonators, Nat. Photonics **5**, 758 (2011).

[21] M. S. Kang, A. Butsch, and P. S. J. Russell, Reconfigurable light-driven opto-acoustic isolators in photonic crystal fibre, Nat. Photonics **5**, 549 (2011).

[22] M. Hafezi and P. Rabl, Optomechanically induced non-reciprocity in microring resonators, Opt. Express **20**, 7672 (2012).

[23] A. Metelmann and A. A. Clerk, Nonreciprocal photon transmission and amplification via reservoir engineering, Phys. Rev. X **5**, 021025 (2015).

[24] X.-W. Xu and Y. Li, Optical nonreciprocity and optomechanical circulator in three-mode optomechanical systems, Phys. Rev. A **91**, 053854 (2015).

[25] Z. Shen, Y.-L. Zhang, Y. Chen, C.-L. Zou, Y.-F. Xiao, X.-B. Zou, F.-W. Sun, G.-C. Guo, and C.-H. Dong, Experimental realization of optomechanically induced non-reciprocity, Nat. Photonics **10**, 657 (2016).

[26] F. Ruesink, M.-A. Miri, A. Alù, and E. Verhagen, Non-reciprocity and magnetic-free isolation based on optomechanical interactions, Nat. Commun. **7**, 13662 (2016).

[27] K. Fang, J. Luo, A. Metelmann, M. H. Matheny, F. Marquardt, A. A. Clerk, and O. Painter, Generalized non-reciprocity in an optomechanical circuit via synthetic magnetism and reservoir engineering, Nat. Phys. **13**, 465 (2017).

[28] L. D. Toth, N. R. Bernier, A. Nunnenkamp, A. K. Feofanov, and T. J. Kippenberg, A dissipative quantum reservoir for microwave light using a mechanical oscillator, Nat. Phys. **13**, 787 (2017).

[29] S. Barzanjeh, M. Wulf, M. Peruzzo, M. Kalaee, P. B. Dieterle, O. Painter, and J. M. Fink, Mechanical on-chip microwave circulator, Nat. Commun. **8**, 953 (2017).

[30] S. Barzanjeh, M. Aquilina, and A. Xuereb, Manipulating the flow of thermal noise in quantum devices, Phys. Rev. Lett. **120**, 060601 (2018).

[31] F. Campaioli, S. Gherardini, J. Q. Quach, M. Polini, and G. M. Andolina, Colloquium: Quantum batteries, arXiv:2308.02277.

[32] R. Alicki and M. Fannes, Entanglement boost for extractable work from ensembles of quantum batteries, Phys. Rev. E **87**, 042123 (2013).

[33] F. Campaioli, F. A. Pollock, F. C. Binder, L. Céleri, J. Goold, S. Vinjanampathy, and K. Modi, Enhancing the charging power of quantum batteries, Phys. Rev. Lett. **118**, 150601 (2017).

[34] J.-Y. Gyhm, D. Šafránek, and D. Rosa, Quantum charging advantage cannot be extensive without global operations, Phys. Rev. Lett. **128**, 140501 (2022).

[35] G. M. Andolina, M. Keck, A. Mari, V. Giovannetti, and M. Polini, Quantum versus classical many-body batteries, Phys. Rev. B **99**, 205437 (2019).

[36] F. C. Binder, S. Vinjanampathy, K. Modi, and J. Goold, Quantacell: Powerful charging of quantum batteries, New J. Phys. **17**, 075015 (2015).

[37] D. Farina, G. M. Andolina, A. Mari, M. Polini, and V. Giovannetti, Charger-mediated energy transfer for quantum batteries: An open-system approach, Phys. Rev. B **99**, 035421 (2019).

[38] K. V. Hovhannisyan, M. Perarnau-Llobet, M. Huber, and A. Acín, Entanglement generation is not necessary for optimal work extraction, Phys. Rev. Lett. **111**, 240401 (2013).

[39] D. Ferraro, M. Campisi, G. M. Andolina, V. Pellegrini, and M. Polini, High-power collective charging of a solid-state quantum battery, Phys. Rev. Lett. **120**, 117702 (2018).

[40] J. Q. Quach, K. E. McGhee, L. Ganzer, D. M. Rouse, B. W. Lovett, E. M. Gauger, J. Keeling, G. Cerullo, D. G. Lidzey, and T. Virgili, Superabsorption in an organic microcavity: Toward a quantum battery, Sci. Adv. **8**, 3160 (2022).

[41] F. Barra, Dissipative charging of a quantum battery, Phys. Rev. Lett. **122**, 210601 (2019).

[42] S. Bhattacharjee and A. Dutta, Quantum thermal machines and batteries, Eur. Phys. J. B **94**, 239 (2021).

[43] J. Q. Quach and W. J. Munro, Using dark states to charge and stabilize open quantum batteries, Phys. Rev. Appl. **14**, 024092 (2020).

[44] X. Yang, Y.-H. Yang, M. Alimuddin, R. Salvia, S.-M. Fei, L.-M. Zhao, S. Nimmrichter, and M.-X. Luo, Battery capacity of energy-storing quantum systems, Phys. Rev. Lett. **131**, 030402 (2023).

[45] J.-Y. Gyhm and U. R. Fischer, Beneficial and detrimental entanglement for quantum battery charging, AVS Quantum Sci. **6**, 012001 (2024).

[46] G. M. Andolina, M. Keck, A. Mari, M. Campisi, V. Giovannetti, and M. Polini, Extractable work, the role of correlations, and asymptotic freedom in quantum batteries, Phys. Rev. Lett. **122**, 047702 (2019).

[47] R. Salvia, M. Perarnau-Llobet, G. Haack, N. Brunner, and S. Nimmrichter, Quantum advantage in charging cavity and spin batteries by repeated interactions, Phys. Rev. Res. **5**, 013155 (2023).

[48] R. R. Rodríguez, B. Ahmadi, G. Suárez, P. Mazurek, S. Barzanjeh, and P. Horodecki, Optimal quantum control of charging quantum batteries, New J. Phys. **26**, 043004 (2024).

[49] F. Mazzoncini, V. Cavina, G. M. Andolina, P. A. Erdman, and V. Giovannetti, Optimal control methods for quantum batteries, Phys. Rev. A **107**, 032218 (2023).

[50] R. R. Rodríguez, B. Ahmadi, P. Mazurek, S. Barzanjeh, R. Alicki, and P. Horodecki, Catalysis in charging quantum batteries, Phys. Rev. A **107**, 042419 (2023).

[51] C.-K. Hu, J. Qiu, P. J. P. Souza, J. Yuan, Y. Zhou, L. Zhang, J. Chu, X. Pan, L. Hu, J. Li, Y. Xu, Y. Zhong, S. Liu, F. Yan, D. Tan, R. Bachelard, C. J. Villas-Boas, A. C. Santos, and D. Yu, Optimal charging of a superconducting quantum battery, Quantum Sci. Technol. **7**, 045018 (2022).

[52] J. Joshi and T. S. Mahesh, Experimental investigation of a quantum battery using star-topology NMR spin systems, Phys. Rev. A **106**, 042601 (2022).

[53] J. Quach, G. Cerullo, and T. Virgili, Quantum batteries: The future of energy storage?, Joule **7**, 2195 (2023).

[54] J. F. Poyatos, J. I. Cirac, and P. Zoller, Quantum reservoir engineering with laser cooled trapped ions, Phys. Rev. Lett. **77**, 4728 (1996).

[55] A. Metelmann and A. A. Clerk, Quantum-limited amplification via reservoir engineering, Phys. Rev. Lett. **112**, 133904 (2014).

[56] A. Metelmann and A. A. Clerk, Nonreciprocal photon transmission and amplification via reservoir engineering, Phys. Rev. X **5**, 021025 (2015).

[57] J. Kerckhoff, K. Lalumière, B. J. Chapman, A. Blais, and K. W. Lehnert, On-chip superconducting microwave circulator from synthetic rotation, Phys. Rev. Appl. **4**, 034002 (2015).

[58] B. J. Chapman, E. I. Rosenthal, J. Kerckhoff, B. A. Moores, L. R. Vale, J. A. B. Mates, G. C. Hilton, K. Lalumière, A. Blais, and K. W. Lehnert, Widely tunable on-chip microwave circulator for superconducting quantum circuits, Phys. Rev. X **7**, 041043 (2017).

[59] Y.-P. Wang, J. W. Rao, Y. Yang, P.-C. Xu, Y. S. Gui, B. M. Yao, J. Q. You, and C.-M. Hu, Nonreciprocity and unidirectional invisibility in cavity magnonics, Phys. Rev. Lett. **123**, 127202 (2019).

[60] M. Kim, A. Tabesh, T. Zegray, S. Barzanjeh, and C.-M. Hu, Nonreciprocity in cavity magnonics at millikelvin temperature, J. Appl. Phys. **135** (2024).

[61] See Supplemental Materials at http://link.aps.org/supplemental/10.1103/PhysRevLett.132.210402 for additional information on calculating stationary energy, optimization procedure for asymmetric damping rates, and studying the advantages of nonreciprocity across different time intervals.

[62] G. M. Andolina, D. Farina, A. Mari, V. Pellegrini, V. Giovannetti, and M. Polini, Charger-mediated energy transfer in exactly solvable models for quantum batteries, Phys. Rev. B **98**, 205423 (2018).

[63] Y.-X. Wang, C. Wang, and A. A. Clerk, Quantum nonreciprocal interactions via dissipative gauge symmetry, PRX Quantum **4**, 010306 (2023).

[64] A. Kossakowski, On quantum statistical mechanics of non-Hamiltonian systems, Rep. Math. Phys. **3**, 247 (1972).

[65] G. Chiribella, Y. Yang, and R. Renner, Fundamental energy requirement of reversible quantum operations, Phys. Rev. X **11**, 021014 (2021).

[66] H. Tajima, N. Shiraishi, and K. Saito, Uncertainty relations in implementation of unitary operations, Phys. Rev. Lett. **121**, 110403 (2018).

[67] H. Tajima, N. Shiraishi, and K. Saito, Coherence cost for violating conservation laws, Phys. Rev. Res. **2**, 043374 (2020).