Letter

Global-local duality of energetic control cost in multipartite quantum correlated systems

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Multipartite quantum correlated systems (MQCSs) are widely utilized in diverse quantum information tasks, where their sophisticated control inherently incurs energetic costs. However, the fundamental characteristics of these control costs remain elusive, largely due to the lack of thermodynamic descriptions capable of capturing the full complexities of MQCSs. Here, we uncover universal thermodynamic relations for arbitrary MQCSs weakly coupled to a thermal bath, establishing an intrinsic global-local duality of control costs. Using these relations, we elucidate the exact role of multipartite correlation—a defining quantum feature of MQCSs—in shaping control costs at finite times. We also demonstrate that the relative magnitude between global and local control costs is undetermined, which perplexes the cost management of MQCSs under finite-time controls. Our results are numerically corroborated with applications to experimentally realizable multiqubit systems undergoing finite-time qubit reset processes.

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Introduction. Realizing quantum technologies that outperform their classical counterparts requires careful energetic considerations [1,2]. Since quantum tasks are generally carried out by dynamical quantum systems, they proceed along with changes in the energetic and entropic characteristics of these systems. Portraying this intertwined feature of quantum tasks necessitates an interdisciplinary perspective, exemplified by the pursuit of connections between information and energy for understanding the intersection between information science and thermodynamics, as pioneered by Maxwell, Szilard, and Landauer and proceeded with many others [3–8]. Indeed, the advancement of quantum technologies can benefit from this interdisciplinary perspective, which not only reveals the constraints on the performance of quantum tasks imposed by fundamental thermodynamic laws [4–10], but also facilitates the study of the emerging issue of their associated thermodynamic cost [1,2,11–19].

Multipartite quantum correlated systems (MQCSs), which hold significant promise for quantum information applications [20–36], provide compelling platforms to integrate the principles of quantum information science and quantum thermodynamics. In this regard, different exotic aspects pertaining to quantum correlation in quantum thermodynamics have been reported in the literature which greatly expand the structure of quantum thermodynamics; notable examples include refined thermodynamic relations [37,38], correlation-involved resource theory [39], correlation-assisted work extraction [40–42], and correlation-induced anomalous energy transfer [43,44] to name just a few. In contrast, the energetic cost of controlling MQCSs—termed control cost from here on—remains poorly understood, despite the recognition that MQCSs, when operated with well-designed control protocols, can advance the potential of various quantum information tasks [20,28–36]. Existing studies in this direction have largely focused on either closed systems [45] or simplified open bipartite settings under additional assumptions [14,46], leaving a general understanding of thermodynamic control cost in generic MQCSs still lacking.

We here focus on thermodynamic control cost of generic MQCSs that are immersed in a thermal bath and subjected to arbitrary finite-time control protocols (see Fig. 1 for a sketch). By combining global and local thermodynamic descriptions, we uncover fundamental relations that embody a general global-local duality of the control cost in MQCSs, requiring the minimum assumption of weak system-bath couplings. Specifically, we reveal an exact sum rule [Eq. (2)] for the globally evaluated control cost, which encompasses contributions from quantum multipartite mutual information [25,45,47,48]—quantifying multipartite total correlation of MQCSs—and dissipated work [49,50], which signifies finite-time thermodynamic processes. We thus add a previously unidentified structure to the collection of energy-information links [3–8]. We also identify a general lower bound [Eq. (3)]

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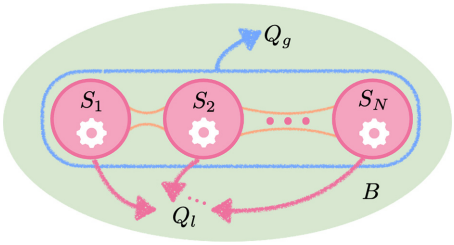


FIG. 1. Schematic picture of the study. We consider MQCSs with several local parties $\{S_i\}$ that are weakly coupled to a thermal bath (B) and subjected to arbitrary control protocols. By examining both the global cost Q_g and the local one Q_l, evaluated at the global and local levels, respectively, we unveil general relations [cf. Eqs. (2) and (3)] that reflect a general global-local duality in the control cost.

on the contrast between globally and locally evaluated costs (Fig. 1) in terms of global dissipated work and changes in intrasystem interaction energy.

Using the established relations, we elucidate the exact role of multipartite correlation in shaping control costs. Notably, the interplay between multipartite correlation and dissipated work leads to an indefinite relative magnitude between globally and locally evaluated control costs, introducing significant complexity in the thermodynamic cost management of MQCSs under finite-time control protocols. We substantiate the theoretical findings using experimentally feasible multiqubit systems undergoing a finite-time qubit reset process. The thermodynamic framework presented here allows the description of MQCSs with arbitrary complexity, an aspect crucial for describing current experimental platforms [20,28–36] and advancing the thermodynamics of control in complex quantum systems.

General setup and main results. We consider a generic MQCS that is weakly coupled to a thermal bath at a temperature T and experiences time-dependent control protocols. The Hamiltonian of the MQCS reads (setting $\hbar = 1$ and $k_B = 1$ hereafter)

$$H_g(t) = \sum_i H_i(t) + H_{\rm I}(t). \tag{1}$$

Here, $\{H_i(t)\}$ denotes the set of Hamiltonian for the driven local parties of the global system. $H_{\rm I}(t)$ encompasses all intrasystem interaction terms, with coupling strengths that may assume any magnitudes depending on the concrete setups. Unlike previous theoretical studies that often assumed noninteracting local parties for simplicity, we consider $H_{\rm I}(t) \neq 0$ to align with experimental conditions. In scenarios where intrasystem couplings cannot be controlled, $H_{\rm I}$ becomes time independent.

The setup in Eq. (1) not only represents a broad class of realistic many-body physical systems, but also captures processes such as quantum measurement and feedback control as special cases where effective bipartite descriptions are applicable [51–53]. We do not include the Hamiltonian of the bath and system-bath interaction in Eq. (1), as they are irrelevant for getting the following main results. The evolution of the global system is described by a time-evolved density matrix $\rho_g(t)$, with $\rho_i(t) = {\rm Tr}_{\bar{i}}[\rho_g(t)]$ being the marginal for the ith local party after tracing out degrees of freedom $\bar{i}$ complementing that of the subsystem i.

We decompose the change in internal energy of the global system $E_g(t) \equiv {\rm Tr}[H_g(t)\rho_g(t)]$ according to the first law of thermodynamics, $\Delta E_g(t) = W_g(t) - Q_g(t)$. Here, $W_g(t) \equiv \int_0^t {\rm Tr}[\rho_g(\tau)\frac{d}{d\tau}H_g(\tau)]d\tau$ denotes the amount of work performed on the global system by external driving fields during the time interval $[0, t]$ [54–56], and $Q_g(t)$ represents the heat dissipation flowing out of the global system which we refer to as the global control cost. For local parties, we can also decompose the energy change $\Delta E_i(t) = W_i(t) - Q_i(t)$ with $E_i(t) = {\rm Tr}[H_i(t)\rho_i(t)]$ and $W_i(t) \equiv \int_0^t {\rm Tr}[\rho_i(\tau)\frac{d}{d\tau}H_i(\tau)]d\tau$ once local information $\{H_i(t), \rho_i(t)\}$ is available. We refer to $Q_l(t) = \sum_i Q_i(t)$ as the local control cost as it permits a local access. We emphasize that the global and local control costs are distinguished by their reliance on global and local measurements, respectively, with the control fields in $H_g(t)$ remaining the same during evaluations.

We find that the global and local control costs are connected through an exact sum rule,

$$Q_g(t) = Q_l(t) + T\,\Delta I_g(t) + W_{\rm dis}^{\Delta}(t). \tag{2}$$

The proof is postponed to the end of the Letter. Here, $\Delta I_g(t) = I_g(t) - I_g(0)$ denotes the change in quantum multipartite mutual information $I_g(t) \equiv \sum_i S_i(t) - S_g(t)$ [25,45,48] during the process. $I_g(t)$ quantifies the multipartite total correlation within the MQCS with $S_{i,g}(t) = -{\rm Tr}[\rho_{i,g}(t)\ln\rho_{i,g}(t)]$ the von Neumann entropy, and reduces to the usual quantum mutual information in bipartite settings. We have defined $W_{\rm dis}^{\Delta}(t) \equiv W_{\rm dis}^{g}(t) - W_{\rm dis}^{l}(t)$ to denote the contrast between global and local dissipated work [49,50] which are, respectively, defined as $W_{\rm dis}^{g}(t) \equiv W_g(t) - \Delta F_g(t)$ and $W_{\rm dis}^{l}(t) \equiv \sum_i [W_i(t) - \Delta F_i(t)]$ with $\Delta F_j(t)$ $(j = i, g)$ denoting the change in nonequilibrium free energy $F_j(t) = E_j(t) - TS_j(t)$ [6,57,58]. Notably, Eq. (2) represents an energy-information link, strengthening the synergy between quantum thermodynamics and quantum information science in MQCSs.

The sum rule Eq. (2) indicates that the global and local costs can differ in MQCSs. Notably, we find that the second law of thermodynamics, when applied to the global system, imposes a general lower bound on the magnitude of the cost contrast $Q_g(t) - Q_l(t)$,

$$Q_g(t) - Q_l(t) \geqslant -W_{\rm dis}^{g}(t) - \Delta E_{\rm I}(t). \tag{3}$$

Here, $\Delta E_{\rm I}(t)$ marks the change in the interaction energy $E_{\rm I}(t) = {\rm Tr}[H_{\rm I}(t)\rho_g(t)]$. For clarity, we provide the proof details at the end of the Letter. The equality of Eq. (3) is attained in the thermodynamic reversible limit. Hence, we expect that the lower bound becomes loose for finite-time processes and cannot be utilized to infer the exact magnitude of the cost contrast.

Equations (2) and (3) constitute our main results. We emphasize that Eqs. (2) and (3) are derived without imposing assumptions on the details of the Hamiltonian in Eq. (1). Hence, they establish a general global-local duality of control costs in MQCSs, which holds across the finite-time regime to the thermodynamic reversible limit. In the special limit of $H_{\rm I}(t) = 0$, we can analytically demonstrate that $Q_g(t) = Q_l(t)$, regardless of the magnitude of multipartite correlation;

We relegate proof details to the Supplemental Material [59]. In this case, the lower bound in Eq. (3) reduces to the maximum work principle for the global system, $W^g_{\rm dis}(t) \geqslant 0$. However, we emphasize that $H_{\rm I}(t) \neq 0$ represents the typical scenario in experiments [28–36], rendering Eqs. (2) and (3) generally nontrivial universal relations without existing counterparts.

It is worth delineating critical properties and implications of Eqs. (2) and (3). (i) Equation (2) indicates that the interplay between multipartite correlation and dissipated work renders the cost contrast $Q_g(t) - Q_l(t)$, while bounded from below as stated in Eq. (3), lacking a definite sign. This perplexes the cost management of MQCSs under finite-time operations. (ii) By examining the asymptotic behavior of Eq. (2) in the thermodynamic reversible limit where $W^g_{\rm dis}(t)$ vanishes, we point out that the dissipated work contrast $W^\Delta_{\rm dis}(t)$ remains nonzero. This occurs because the local parties of a correlated system no longer adhere to the standard form of the second law of thermodynamics [38], as demonstrated in [59]. (iii) We remark that the global cost $Q_g(t)$ should comply with the standard form of the Landauer principle, $Q_g(t) \geqslant -T\Delta S_g(t)$ [9]. In comparison, by combining this inequality with Eq. (2), we find that the local cost satisfies a modified form of the Landauer principle, $Q_l(t) \geqslant -T\sum_i \Delta S_i(t) - W^\Delta_{\rm dis}(t) \neq -T\sum_i \Delta S_i(t)$, in light of the second point mentioned above. (iv) Finally, we note that maintaining multipartite correlation in MQCSs over time is essential for unlocking their scalability advantages. In this regard, Eqs. (2) and (3) translate into thermodynamic constraints on the magnitude of $\Delta I_g(t)$ that can be actively achieved in MQCSs. Specifically, Eq. (2) implies that tuning $\Delta I_g(t)$ necessarily induces global-local contrasts in control cost and dissipated work. Combining Eqs. (2) and (3), we find that $\Delta I_g(t)$ is subject to a thermodynamic lower bound, $T\Delta I_g(t) \geqslant W^l_{\rm dis}(t) - 2W^g_{\rm dis}(t) - \Delta E_{\rm I}(t)$.

Numerical demonstration. We validate our main theoretical findings using multiqubit systems weakly coupled to a thermal bath at a temperature $T = \beta^{-1}$ with open boundaries,

$$H_g(t) = \frac{\varepsilon_t}{2}\sum_{i=1}^{N}\left[\cos(\theta_t)\sigma_i^z + \sin(\theta_t)\sigma_i^x\right] + \sum_{i=1}^{N-1}\lambda_t\sigma_i^x\sigma_{i+1}^x. \tag{4}$$

Here, N marks the number of qubits, $\{\sigma_i^{x,z}\}$ represent the Pauli matrices, and ε_t, θ_t, and λ_t are subject to time-dependent control protocols. The multiqubit model with nearest-neighbor couplings generalizes a single-qubit setup [58,60,61], and can be experimentally realized using coupled quantum dot systems [62] and tight-binding superconducting qubit arrays [63], where both the local parties and couplings are tunable. We here consider the following finite-time control fields $\varepsilon_t = \varepsilon_0 + (\varepsilon_\tau - \varepsilon_0)\sin(\pi t/2\tau)^2$, $\theta_t = \frac{\pi}{2}(t/\tau - 1)$ with τ setting the operation time limit and $\lambda_t = \lambda\cos(\theta_t)$. We take these experimentally feasible control fields [58,60,61] to implement a multiqubit reset process, a procedure with significant applications in numerous quantum algorithms [64] and quantum error correction [65,66].

At weak couplings, the evolution of $\rho_g(t)$ is governed by the quantum Lindblad master equation

$$\frac{\rm d}{{\rm d}t}\rho_g(t) = -i[H_g(t), \rho_g(t)] + \sum_{\mu=1,2}\gamma_\mu\mathcal{D}[L_\mu]\rho_g(t). \tag{5}$$

Here, $\gamma_\mu \geqslant 0$ is the damping coefficient of quantum channel μ, and $\mathcal{D}[L_\mu]\rho = L_\mu\rho L_\mu^\dagger - \frac{1}{2}\{L_\mu^\dagger L_\mu, \rho\}$ is the Lindblad superoperator with L_μ being the Lindblad jump operator and $\{A, B\} = AB + BA$. For the present model, the Lindblad jump operators, due to the coupling to a common thermal bath, read [67] $L_1(t) = \sqrt{\varepsilon(t)[N_B(t)+1]}\sum_i\sigma_i^-$ and $L_2(t) = \sqrt{\varepsilon(t)[N_B(t)]}\sum_i\sigma_i^+$, where $\sigma_i^{+,-}$ denote spin ladder operators of the ith qubit and $N_B(t) = [e^{\beta\varepsilon(t)} - 1]^{-1}$. For the sake of simplicity, we take $\gamma_{1,2} = \gamma$ in simulations. A time step $\delta t = 1 \times 10^{-4}$ is used to numerically integrate Eq. (5).

For the demonstration purpose, we choose an initial thermal state $\rho_g(0) = e^{-\beta H_g(0)}/{\rm Tr}[e^{\beta H_g(0)}]$ for the global system and depict a typical set of numerical results in Fig. 2. In Figs. 2(a) and 2(b) for $N = 2$ and $N = 4$, respectively, we show numerical results for quantities on two sides of Eq. (2). The perfect agreement between both clearly demonstrates the validity of the sum rule in multiqubit systems. From the result, we note that the global cost $Q_g(t)$ monotonically increases as the system evolves. In Figs. 2(c) and 2(d), we verify Eq. (3) by presenting numerical results for the cost contrast $Q_g(t) - Q_l(t)$ and its lower bound (LB) as given by the right-hand-side of Eq. (3). We observe that $Q_g(t) - Q_l(t) > 0$, especially at large times, from both Figs. 2(c) and 2(d).

To analyze the sign of $Q_g(t) - Q_l(t)$ in the present model, we rewrite the sum rule as $Q_g(t) - Q_l(t) = T\Delta I_g(t) + W^\Delta(t) - F^\Delta(t)$, where we denote $F^\Delta(t) \equiv \Delta F_g(t) - \sum_i \Delta F_i(t)$ and $W^\Delta(t) \equiv W_g(t) - \sum_i W_i(t)$. Specifically, $F^\Delta(t)$ represents the contrast in work capacity (i.e., free energy change) between the global system and its local parties. Since intrasystem correlations act as a thermodynamic resource that enables additional work extraction [41,46], it is natural to expect that the global system has a greater work capacity than its local parties, leading to $F^\Delta(t) > 0$ as confirmed numerically [59]. The term $W^\Delta(t)$ becomes nonzero when the intrasystem interaction is actively controlled. In our model, we continuously increase the strength of the intrasystem interaction, requiring an additional work performed on the system such that $W^\Delta(t) > 0$ [59]. Additionally, we observe $\Delta I_g(t) > 0$ from the insets of Figs. 2(a) and 2(b), indicating that control fields generate multipartite correlations. Since creating correlations inherently incurs an energetic cost as exactly quantified by $T\Delta I_g(t)$ [45], we can attribute the sign of $Q_g(t) - Q_l(t)$ here to the interplay between the gain in work capacity $F^\Delta(t)$ and energetic costs of building up correlation and interaction as captured by $T\Delta I_g(t) + W^\Delta(t)$.

For the scenario considered in Fig. 2, we find that the costs outweigh the gain, leading to a positive cost contrast $Q_g(t) - Q_l(t) > 0$. However, when the gain instead surpasses the costs, the scenario of $Q_g(t) - Q_l(t) < 0$ can also occur, as demonstrated using a multiqubit model with a modified intrasystem interaction and the same control fields (see details in [59]). The varying signs of $Q_g(t) - Q_l(t)$ in multiqubit systems confirm our expectations that the cost contrast lacks a definite sign, and that the sum rule captures the rich dynamical behaviors of control costs in MQCSs.

Proof of Eqs. (2) and (3). To proof Eq. (2), we start from the nonequilibrium free energy [6,57,58] of the global system, $F_g(t) = E_g(t) - TS_g(t)$. Combining the first law of thermodynamics for the global system and the

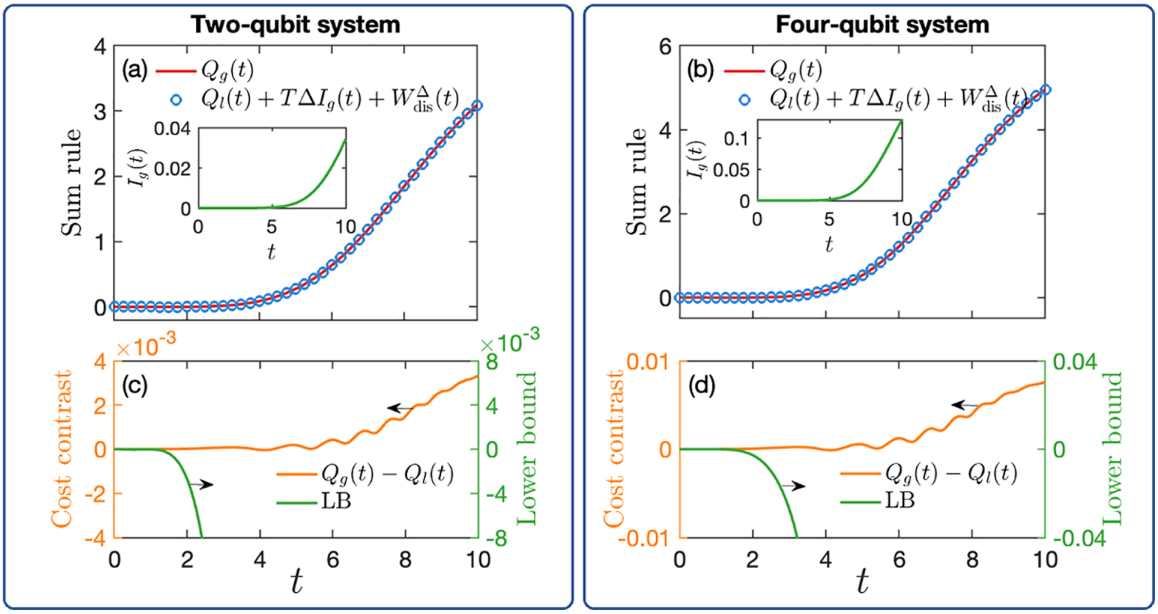


FIG. 2. Examining Eqs. (2) and (3) in multiqubit systems with $N = 2$ (left panel) and $N = 4$ (right panel). (a), (b) Validity of the sum rule for the global cost $Q_g(t)$. (c), (d) The cost contrast $Q_g(t) - Q_l(t)$ (solid line, left axis) and its lower bound (LB) (dashed line, right axis). Insets: Time-dependent quantum multipartite mutual information $I_g(t)$. Other parameters are $\beta = 1$, $\lambda = 0.02$, $\gamma = 0.02$, $\varepsilon_0 = 0.4$, $\varepsilon_\tau = 10$, and $\tau = 10$.

definition of quantum multipartite mutual information [25,45,48], the change in $F_g(t)$ can be expressed as $\Delta F_g(t) = W_g(t) - Q_g(t) - T[\sum_i \Delta S_i(t) - \Delta I_g(t)]$. To proceed, we further consider nonequilibrium free energies $F_i(t) = E_i(t) - TS_i(t)$ for local parties which can be adopted to replace the term $\sum_i \Delta S_i(t)$ in $\Delta F_g(t)$, yielding the following relation:

$$Q_g(t) = W_g(t) - \Delta F_g(t) + T\Delta I_g(t) + \sum_i [\Delta F_i(t) - \Delta E_i(t)]. \tag{6}$$

The sum rule Eq. (2) can be immediately derived from Eq. (6) by invoking the energy decomposition for local parties $\Delta E_i(t) = W_i(t) - Q_i(t)$ and applying the definition of dissipated work $W_{\rm dis} \equiv W - \Delta F$ [49,50] to both the global system and local parties.

We then turn to the cost contrast $Q_g(t) - Q_l(t)$. Inserting the energy decomposition $E_g(t) = \sum_i E_i(t) + E_{\rm I}(t)$ into the nonequilibrium free energy $F_g(t) = E_g(t) - TS_g(t)$ of the global system and considering the latter's change, we get

$$T\Delta S_g(t) = \sum_i \Delta E_i(t) + \Delta E_{\rm I}(t) - \Delta F_g(t). \tag{7}$$

We next invoke the relation $\Delta E_i(t) = W_i(t) - Q_i(t)$ and transform Eq. (7) into

$$T\Delta S_g(t) = -Q_l(t) + \Delta E_{\rm I}(t) + W^g_{\rm dis}(t). \tag{8}$$

To arrive at the above form, we have utilized the fact that $W_g(t) = \sum_i W_i(t)$ since only the local system Hamiltonian $\{H_i(t)\}$ is being driven.

To proceed, we further add $Q_g(t)$ to both sides of Eq. (8) and utilize the second law of thermodynamics satisfied by the global system $T\Delta S_g(t) + Q_g(t) \geqslant 0$, yielding

$$Q_g(t) - Q_l(t) + \Delta E_{\rm I}(t) + W^g_{\rm dis}(t) \geqslant 0. \tag{9}$$

By rearranging the terms, we arrive at Eq. (3). From the above derivations, we clearly see that Eqs. (2) and (3) are general.

Discussion and conclusion. The experimental verifications of Eqs. (2) and (3), though challenging due to the necessity of global measurements on $\rho_g(t)$, is well within the reach of current experimental capabilities. We remark that the detailed form of $H_g(t)$ is known *a priori* in typical experiments [28–36]. $\rho_g(t)$ can be measured using quantum state tomography techniques advanced for many-body systems such as qubit arrays [68–70], whose capability can be further enhanced by incorporating machine learning methods [71–73].

We note that our main results [cf. Eqs. (2) and (3)] require the existence of a fixed bath temperature. However, in realistic nanoscale environments, which are often finite in size, time-dependent effective temperatures provide a more accurate description [74–80]. To accommodate such situations, we can adopt a generalized definition of the nonequilibrium free energy [58] which can incorporate a time-dependent effective temperature. We show that the basic form of the sum rule remains intact, with just the term $T\Delta I_g(t)$ in Eq. (2) replaced by $T(t)I_g(t) - T(0)I_g(0)$. However, the lower bound on $Q_g(t) - Q_l(t)$ as given by Eq. (3) cannot be straightforwardly generalized (see details in [59]). Thus, we expect that at least the global-local duality underpinning the sum rule may represent a general feature of a wide range of nonequilibrium MQCSs, without coupling to a thermal bath.

In summary, we provided a general thermodynamic description for addressing the control cost of arbitrary MQCSs under finite-time operations. We uncover a fundamental global-local duality governing the control cost of MQCSs, elucidating the interplay between relevant thermodynamic and information-theoretic quantities. The critical properties and implications of our results are delineated and demonstrated

using experimentally feasible multiqubit models. We believe that these findings offer fresh insights into the thermodynamics of control in MQCSs, which are particularly relevant for applications of MQCSs in quantum simulation, error correction and beyond.

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Data availability. The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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