Accelerating Quantum Relaxation via Temporary Reset: A Mpemba-Inspired Approach

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Slow relaxation processes spanning widely separated timescales pose fundamental challenges for probing steady-state properties and engineering functional quantum systems, such as quantum heat engines and quantum computing devices. We introduce a protocol that enables significant acceleration of relaxation in general Markovian open quantum systems by temporarily coupling the system to a reset channel, inspired by the Mpemba effect. Crucially, this acceleration persists even when the slowest decaying Lindbladian modes form complex-conjugate pairs. Unlike previous approaches, which typically target a single mode, our protocol may suppress multiple relaxation modes simultaneously. This framework provides a versatile and experimentally feasible tool for controlling relaxation timescales, with broad implications for quantum thermodynamics, computation, and state preparation.

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Introduction—Relaxation processes in open quantum systems, wherein a driven or interacting system settles into an equilibrium or a nonequilibrium steady state, are of fundamental importance in nonequilibrium physics and for practical quantum technologies. The timescales of such relaxations often determine the feasibility and performance of quantum devices. For example, in quantum thermodynamics, the power output of cyclic heat engines or transport junctions can be limited by how rapidly the working substance relaxes to its steady state in each cycle [1,2]. In addition, fast relaxation facilitates ground state laser cooling [3–5] and the reliable preparation of quantum states [6]. Faster relaxation can also lead to more efficient quantum algorithms [7], particularly in dissipative computational architectures where the output is encoded in the system's stationary state [7,8]. Slow relaxation inevitably allows other unwanted dissipative dynamics to consume substantial time resources, compromising the efficiency of the final stabilized state [9,10]. In contrast, slow relaxation is sometimes desirable—such as when the computational task targets metastable states [11–14], whose prolonged lifetimes are essential for robust information storage or approximate optimization.

Researchers have developed various techniques to accelerate relaxation toward both thermal equilibrium [15] and nonequilibrium steady states [10,16,17]. Most of these methods, however, are restricted to specific systems and rely on intricate external controls. For instance, some approaches require auxiliary Hamiltonians or full knowledge of the system Hamiltonian, limiting their applicability to simple setups. A more general and elegant strategy is to accelerate relaxation by tailoring the initial state distribution, without the need for continuous external driving. A particularly intriguing example of this approach is inspired by the so-called Mpemba effect [18], a counterintuitive relaxation anomaly in which a hotter system can cool faster than a colder one. This phenomenon implies the existence of an optimal initial state that minimizes the relaxation timescale. Theoretical frameworks for identifying such optimal states have been proposed in both classical [19–31] and quantum systems [9,32–41]; see also Refs. [42,43] for recent reviews. However, existing quantum implementations typically impose strict constraints, such as assuming pure initial states or requiring the second-largest eigenvalue of the Lindbladian to be real. Moreover, careful initial-state design is delicate since it relies on fine-tuning of control parameters and detailed knowledge of the initial state and system dynamics, which are typically not known *a priori* in complex quantum systems. These issues limit their applicability in general settings. Thus, a general and practical method for inducing faster relaxation in open quantum systems, which is robust with respect to different initial states, remains elusive.

In this Letter, we contribute to addressing these challenges by proposing an experimentally feasible protocol that can both accelerate and decelerate general quantum relaxation processes from any initial state, by resetting the system to a specified (possibly mixed) target state [44–46].

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Here, this type of quantum reset operation can be realized through engineered dissipation, i.e., quantum reservoir engineering [8,47–49], and should be distinguished from the conventional reset typically used for qubit initialization in quantum computation [50–55]. The motivation is more closely related to Refs. [56,57]. In our protocol, we introduce a finite-duration reset phase into the open-system dynamics by temporarily coupling the system to a reset channel. We show that this protocol can significantly accelerate relaxation, even when starting from mixed initial states and in the presence of complex decay modes. Remarkably, our method allows for the simultaneous suppression of multiple relaxation modes—a capability absent in previous approaches. Moreover, selective application of the protocol enables the system of interest to remain in metastable states for extended periods, offering a systematic route to suppress relaxation.

Setup—We consider a general Markovian open quantum system defined on a Hilbert space $\mathcal{H}$ of dimension d, whose dynamics are governed by the Lindblad-Gorini-Kossakowski-Sudarshan master equation $d\rho/dt = \mathcal{L}(\rho)$, where the Lindbladian superoperator $\mathcal{L}$ is given by

$$\mathcal{L}(\rho) \equiv -i[H,\rho] + \sum_i \left[J_i \rho J_i^\dagger - \frac{1}{2}\{J_i^\dagger J_i, \rho\}\right]. \quad (1)$$

Here, H is the Hamiltonian of the system and the jump operators J_i describe the dissipative coupling to the environment. The evolution can be analyzed via the spectral decomposition of $\mathcal{L}$, whose eigenvalues are denoted λ_k and ordered such that $0 \geq \mathrm{Re}(\lambda_k) \geq \mathrm{Re}(\lambda_{k+1})$. Assuming that $\mathcal{L}$ is diagonalizable, let R_k and L_k denote the corresponding right and left eigenmatrices, respectively, satisfying $\mathcal{L}(R_k) = \lambda_k R_k$ and $\mathcal{L}^\dagger(L_k) = \lambda_k^\star L_k, k = 1, \ldots, d^2$, with biorthogonal normalization [58]

$$\mathrm{Tr}(L_k^\dagger R_h) = \delta_{kh}. \quad (2)$$

The dual superoperator $\mathcal{L}^\dagger$ governs the Heisenberg-picture dynamics of observables:

$$\mathcal{L}^\dagger(O) = i[H,O] + \sum_i \left[J_i^\dagger O J_i - \frac{1}{2}\{J_i^\dagger J_i, O\}\right].$$

Given an initial state ρ_0, the system state at time t evolves as

$$\rho(t) = e^{t\mathcal{L}}[\rho_0] = \rho_{\mathrm{ss}} + \sum_{k=2}^{d^2} c_k e^{\lambda_k t} R_k, \quad (3)$$

where $c_k \equiv \mathrm{Tr}(L_k^\dagger \rho_0)$. The unique stationary state of the open quantum system ρ_{ss} is given by $\rho_{\mathrm{ss}} = \lim_{t\to\infty} \rho(t) = R_1$, assuming $\lambda_1 = 0$ is nondegenerate. In the long-time limit, the relaxation is dominated by the slowest decaying mode, and the deviation from stationarity obeys $||\rho(t) - \rho_{\mathrm{ss}}|| \sim \exp(-|\mathrm{Re}\lambda_2|t)$.

Reset protocol to accelerate or decelerate relaxation processes—Recognizing that the slowest decaying mode governs the relaxation timescale, we propose a quantum reset protocol that can suppress or promote its excitation from a general initial state via a finite-duration reset phase, thereby enabling exponential acceleration of relaxation. Specifically, we let the system evolve under modified dynamics for a finite time t_s, during which it is stochastically reset to a chosen state ρ_δ at random times, with events following a Poisson process of rate r. After time t_s, the reset channel is turned off, and the system resumes evolving under the original Lindbladian $\mathcal{L}$, that is, the system density matrix $\rho(t)$ evolves under the reset channel for $t \in [0, t_s]$ and decouples from the channel for $t > t_s$. The modified dynamics during the reset phase $t \in [0, t_s]$ is governed by a new Lindbladian $\mathcal{L}_r$ [44], defined as

$$\mathcal{L}_r(\rho) \coloneqq \mathcal{L}(\rho) + \mathcal{D}_r(\rho) = \mathcal{L}(\rho) + r\mathrm{Tr}(\rho)\rho_\delta - r\rho, \quad (4)$$

where $\mathcal{D}_r(\rho) = r[\mathrm{Tr}(\rho)\rho_\delta - \rho]$. Expressing $\rho_\delta = \sum_\alpha p_\alpha |\psi_\alpha\rangle \times \langle\psi_\alpha|$, $\mathcal{L}_r$ is equivalent to introducing jump operators $J_{i,\alpha}^r = \sqrt{r p_\alpha}|\psi_\alpha\rangle\langle\phi_i|$ into $\mathcal{L}$, where $|\phi_i\rangle$ form an auxiliary complete orthonormal basis (see End Matter). The term $r\mathrm{Tr}(\rho)\rho_\delta$ contributes only to the stationary mode and vanishes on all traceless eigenmodes R_i with $i \geq 2$, which satisfy $\mathrm{Tr}(R_i) = 0$ due to $\mathrm{Tr}(L_1^\dagger R_i) = 0$ and $L_1^\dagger = \mathbb{I}$ [cf. Eq. (2)]. Consequently, the modified Lindbladian acts as

$$\mathcal{L}_r(R_i) = (\lambda_i - r)R_i, i \in \{2, \ldots, d^2\}. \quad (5)$$

This induces a uniform spectral shift by $-r$ for all nonstationary eigenmodes, leaving the eigenmatrices (R_i, $i \geq 2$) themselves unchanged. In contrast, the stationary mode R_1 retains eigenvalue zero but is modified under reset. Let $\rho^r(t) \coloneqq e^{t\mathcal{L}_r}\rho_0$ be the state at time t under reset dynamics. Its spectral decomposition reads as

$$\rho^r(t) = e^{t\mathcal{L}_r}\rho_{\mathrm{ss}} + \sum_{k=2}^{d^2} c_k e^{(\lambda_k - r)t} R_k$$
$$\coloneqq \rho_{\mathrm{ss}} + \sum_k c_k^r(t) e^{\lambda_k t} R_k, (0 \leq t \leq t_s), \quad (6)$$

where $c_k^r(t)$ are time-dependent mode amplitudes. Explicitly, these modified coefficients are given by (End Matter)

$$c_k^r(t) \equiv \left[c_k - \frac{r \cdot d_k}{r - \lambda_k}\right] e^{-rt} + \frac{r \cdot d_k}{r - \lambda_k} e^{-\lambda_k t}, \quad (7)$$

where $d_k \equiv \mathrm{Tr}(L_k^\dagger \rho_\delta)$ is the overlap coefficient of ρ_δ. Notably, Eq. (6) can be generalized to the dynamics of observables (see Supplemental Material [59]), which may be more convenient to measure.

After the reset phase ($t > t_s$), the system returns to the original Lindbladian evolution. The dynamics for $t > t_s$ then reads as

$$\rho^r(t) = \rho_{\rm ss} + \sum_k c_k^r(t_s) e^{\lambda_k t} R_k. \quad (t > t_s). \tag{8}$$

Importantly, the modified spectral decomposition, Eq. (8), takes the same form as the original one, Eq. (3), which allows us to directly compare the relaxation dynamics with and without reset. From this comparison, it is clear that for $t > t_s$ our protocol is equivalent to the original Lindbladian evolution, but with different overlap coefficients,

$$\rho(t) = \rho_{\rm ss} + \sum_{k=2}^{d^2} c'_k e^{\lambda_k t} R_k. \quad (t > t_s). \tag{9}$$

Here, $c'_k \equiv c_k^r(t_s)$ are determined by the reset protocol.

Thus, our protocol controls relaxation dynamics by tuning the overlap coefficients c_k, without requiring any special preparation of the initial state. By choosing an appropriate ρ_δ and varying r and t_s, one can tune the amplitudes of multiple relaxation modes. In particular, ensuring $|c_2^r(t_s)| < |c_2| (> |c_2|)$ suppresses (enhances) the dominant mode, which is sufficient for acceleration (deceleration), ultimately bringing the system closer to (farther from) stationarity [19,32]. Specifically, the dominant mode is fully eliminated and exponential speed-up occurs when $|c_2^r(t_s)| = 0$; this, however, requires fine-tuning and prior knowledge and is therefore challenging to realize in large systems. We thus focus on the weaker but practically relevant scenario $|c_2^r(t_s)| < |c_2|$.

Sufficient conditions for suppression or promotion of relaxation modes—We derive a sufficient condition [59],

$$\mathrm{Re}(c_2^* d_2) < |c_2|^2, \tag{10}$$

under which there always exists a threshold $t_c > 0$ such that, for any $r > 0$ and any $t_s \in (0, t_c]$, one has $|c_2^r(t_s)| < |c_2|$, that is, the relaxation is accelerated for all t_s in this range. If the condition is violated, there exists another threshold $t'_c > 0$ with $|c_2^r(t_s)| \geq |c_2|$ for $t_s \in (0, t'_c]$, leading to deceleration. However, Eq. (10) is not necessary for acceleration since one may still have $|c_2^r(t_s)| < |c_2|$ for $t_s > t'_c$, when it is not satisfied.

Several remarks are in order. First, the condition applies both when λ_2 is real and when it forms a complex-conjugate pair. In the latter case, the long-time dynamics takes the form

$$\rho(t) = \rho_{\rm ss} + e^{\lambda_2 t} c_2 R_2 + e^{\lambda_2^* t} c_2^* R_2^\dagger + \mathcal{O}(e^{\lambda_3 t}), \tag{11}$$

where $c_2 = \mathrm{Tr}(L_2^\dagger \rho_0)$ and $c_2^* = \mathrm{Tr}(L_2 \rho_0)$ is its complex conjugate. Most existing strategies cannot simultaneously suppress both components of such complex relaxation modes. Our protocol, however, naturally addresses this limitation. The modified coefficients under reset are

$$c_2^r(t) = \left[c_2 - \frac{r d_2}{r - \lambda_2}\right] e^{-rt} + \frac{r d_2}{r - \lambda_2} e^{-\lambda_2 t}, \tag{12a}$$

$$c_2^{r,*}(t) = \left[c_2^* - \frac{r d_2^*}{r - \lambda_2^*}\right] e^{-rt} + \frac{r d_2^*}{r - \lambda_2^*} e^{-\lambda_2^* t}. \tag{12b}$$

These coefficients remain complex conjugates for all t, i.e., $[c_2^r(t)]^* = c_2^{r,*}(t)$. Equation (10) then ensures that $|c_2^{r,*}(t_s)| = |c_2^r(t_s)| \leq |c_2| = |c_2^*|$ for a moderate t_s.

Second, Eq. (10) is a weak condition satisfied by many choices of ρ_δ. In particular, it does not require that ρ_δ be closer to the steady state than ρ_0, nor that the overlap coefficient $|d_2|$ be smaller than $|c_2|$. For example, when λ_2 is real, if $\mathrm{Re}(d_2)$ and c_2 have opposite signs, the condition is satisfied even when $|d_2|$ is arbitrarily large. In Example 2 we show that the condition is typically satisfied for a large fraction of randomly chosen initial states.

Third, neither r nor t_s needs to be precisely tuned in advance to achieve acceleration or deceleration. Provided that Eq. (10) holds, any positive r and moderately small t_s are sufficient for acceleration, and t_s can be further optimized *a posteriori* based on the observed dynamics. Our protocol therefore provides a general and robust framework for controlling relaxation in open quantum systems, without the need for fine-tuning.

Finally, it is straightforward to see that

$$\mathrm{Re}(c_k^* d_k) < |c_k|^2 \tag{13}$$

is a sufficient condition for suppression of the kth mode. Since multiple such conditions for different k can be satisfied simultaneously, our protocol allows simultaneous suppression (or promotion) of multiple relaxation modes. In [59] we provide numerical evidence that multiple conditions can indeed be satisfied in parallel, using the model and protocol of Example 2.

So far we have focused on practical scenarios where fine-tuning is not possible and neither the system dynamics nor the initial state are known *a priori*. If these constraints are lifted, as in most previous studies, our protocol can fully eliminate the dominant relaxation mode ($|c_2^r(t_s)| = 0$) by choosing an optimal $t_s = t_s^*$. When λ_2 is real, the exact expression is [24,59]

$$t_s^* = \frac{1}{r - \lambda_2} \ln\left(1 - \frac{c_2(r - \lambda_2)}{r d_2}\right). \tag{14}$$

A sufficient condition ensuring that $t_s^* \geq 0$ is $c_2/d_2 \leq 0$. Further details, including the case of complex-conjugate pairs, are given in [59].

Example 1—To illustrate our protocol, we first consider a minimal example: a two-level quantum system, which may represent a qubit or a spin-1/2 particle. While minimal, this example is crucial in many practical settings, as qubit reset often constitutes the bottleneck in large-scale quantum processes [57,72]. We set the ground state energy $E_0 = 0$ and excited state energy $E_1 = E$, such that the system Hamiltonian reads $H = E|1\rangle\langle 1| + \Omega(\sigma^+ + \sigma^-)$, where Ω is the intrinsic coherent coupling between the two levels, $\sigma^+ = |1\rangle\langle 0|$ and $\sigma^- = |0\rangle\langle 1|$. The system is coupled to a thermal environment through the Lindblad jump operators $J_0 = \sqrt{\gamma_0}\sigma^+$ and $J_1 = \sqrt{\gamma_1}\sigma_-$. Here, γ_0 (γ_1) is the transition rate from state $|0\rangle$ to state $|1\rangle$ (or vice versa). Assuming the rates satisfy detailed balance with inverse temperature $\gamma_0 e^{-\beta_{\rm env}E_0} = \gamma_1 e^{-\beta_{\rm env}E_1}$, i.e., $\gamma_0 = \gamma_1 e^{-\beta_{\rm env}E}$, the system evolves to a unique equilibrium state $\rho_{\rm eq} = [e^{-\beta_{\rm env}H}/\mathrm{Tr}(e^{-\beta_{\rm env}H})]$ irrespective of the initial condition. We initialize the system in a general mixed state:

$$\hat{\rho}_0 = \begin{pmatrix} \frac{1}{1+e^{-\beta_0 E}} & ke^{i\phi} \\ ke^{-i\phi} & \frac{e^{-\beta_0 E}}{1+e^{-\beta_0 E}} \end{pmatrix},$$

$$0 \le k \le \frac{\sqrt{e^{\beta_0 E}}}{e^{\beta_0 E}+1}, \qquad 0 \le \phi \le 2\pi,$$

where the off-diagonal terms $ke^{\pm i\phi}$ quantify initial quantum coherence. When $\Omega = 0$, all eigenmodes and eigenvalues of the Lindbladian can be computed analytically [59]. When we apply our reset protocol with the reset state $\rho_\delta = |0\rangle\langle 0|$, corresponding to the ground state, the two slowest decaying modes form a complex-conjugate pair:

$$c_2^r(t) = e^{-rt+i\phi}k, \qquad c_2^{r,*}(t) = e^{-rt-i\phi}k. \qquad (15)$$

These overlaps are exponentially suppressed compared with the original overlap, i.e., $c_2^r(t) = e^{-rt}c_2$ and $c_2^{r,*}(t) = e^{-rt}c_2^*$. Clearly, the relaxation is accelerated for any r and t_s.

Since any t_s can ensure acceleration, we choose t_s such that a faster decaying mode is eliminated. Specifically, solving $c_4^r(t_s) = 0$ gives such a t_s, whose expression is provided in [59]. For a sufficiently large r, the slowest mode's amplitude becomes exponentially small after the reset phase. Alternative selections for t_s can also achieve accelerated relaxation, indicating that detailed dynamics knowledge is not essential. Notably, if the initial inverse temperature satisfies $\beta_0 < \beta_{\rm env}$, a scenario analogous to "cooling," the ground state $|0\rangle\langle 0|$ serves as a useful ρ_δ regardless of the rate r. Conversely, when $\beta_0 > \beta_{\rm env}$, resetting to the ground state decelerates relaxation, necessitating the choice $\rho_\delta = |1\rangle\langle 1|$ for acceleration. This example demonstrates that, by appropriately choosing ρ_δ, one can switch between reset-induced acceleration and deceleration of relaxation.

We next present numerical results for illustration. We characterize the distance between the transient state $\rho(t)$ and the stationary state $\rho_{\rm ss}$ using the standard trace distance $D[\rho(t)|\rho_{\rm ss}] := \mathrm{Tr}|\rho(t) - \rho_{\rm ss}|/2$, where $|A| := \sqrt{A^\dagger A}$. Other measures, such as the L_∞ norm $D_\infty[\rho(t)|\rho_{\rm ss}] := \max_i|\lambda_i|$ [λ_i is the ith eigenvalue of $\rho(t) - \rho_{\rm ss}$], yield qualitatively similar outcomes. We first plot $D[\rho(t)|\rho_{\rm eq}]$ over t for various rates r in Fig. 1(a). The parameters for $r > 0$ cases are chosen so that they are initially farther from equilibrium compared to the $r = 0$ case (see the caption of Fig. 1). The reset protocol substantially accelerates relaxation so that the former cases reach stationarity faster even if they are farther from stationarity initially. We next plot t_s as a function of r in Figs. 1(b) and 1(d), which show that t_s decreases monotonically as r increases [recall that t_s is chosen so that $c_4^r(t_s) = 0$]. Figure 1(c) confirms accelerated relaxation for $\Omega = 2$, here measured by the L_∞ norm to reduce fluctuations in the curves. Notably, the steady state is not equilibrium when $\Omega > 0$.

Example 2—To illustrate the potential feasibility of our protocol in complex quantum systems, we consider another example, the dissipative transverse-field Ising model (TFIM), a paradigmatic system relevant to various quantum platforms. We numerically study an open dissipative TFIM of length N with Hamiltonian

$$H = -J\sum_{i=1}^{N-1}\sigma_i^z\sigma_{i+1}^z - g\sum_{i=1}^{N}\sigma_i^x \qquad (16)$$

and jump operators $J_{i,\downarrow} = \sqrt{\gamma}\sigma_i^-, J_{i,\uparrow} = \sqrt{\gamma e^{-\beta}}\sigma_i^+$. Here, $\sigma_i^{x,z}$ are Pauli operators, $\sigma_i^- = |0\rangle\langle 1|$ and $\sigma_i^+ = |1\rangle\langle 0|$. We fix $J = 1$, $g = 1.2$, $\gamma = 0.5$, and $\beta = 1/k_BT = 1$ throughout. We choose $\rho_\delta = \mathbb{I}/d$ (with $d = 2^N$). The initial state ρ_0 is chosen by normalizing $\rho_{\rm ss} + \alpha V_2$, with $\rho_{\rm ss}$ being the steady-state and V_2 being the second normalized eigenvector

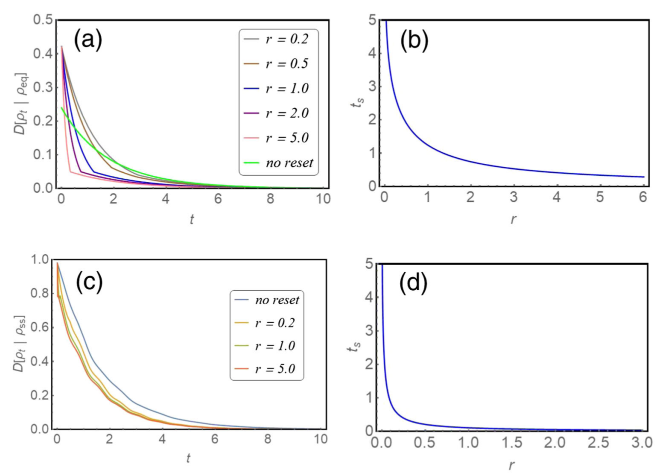

FIG. 1. Relaxation dynamics of the two-state system. (a), (c) Distances between $\rho(t)$ and $\rho_{\rm eq}$ as a function of time t with different resetting rates r. (b),(d) The t_s as a function of the resetting rate r. The parameters are chosen as $\beta_0 = 2.0$, $k = 0.32$, and $\phi = 1$ with reset protocol and $\beta_0 = 3.0$, $k = 0.21$, and $\phi = 1$ without reset. The environment is at lower temperature $\beta_{\rm env} = 4.0$. $\gamma_1 = 1.0$. For (a),(b) $\Omega = 0$, and for (c),(d) $\Omega = 2.0$.

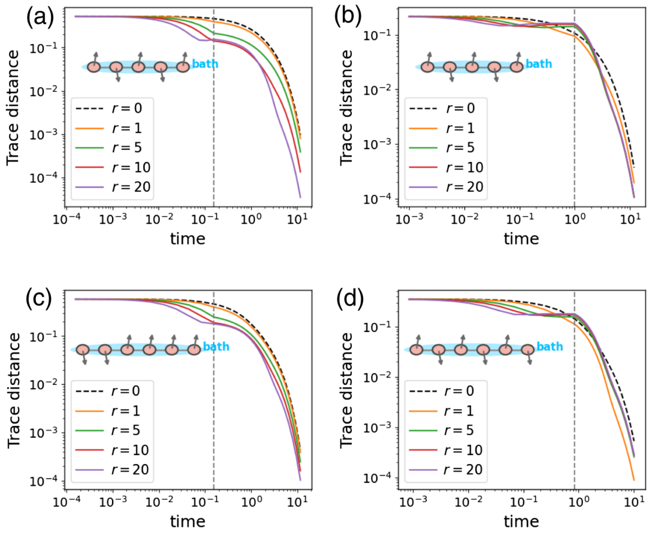

FIG. 2. Acceleration of the relaxation of TFIMs. Reset rates r are chosen as 0, 1.0, 5.0, 10.0, 20.0. (a) $N = 5$ $(d = 32)$, $\alpha = 0.55$, $t_s = 0.08\tau_2$. (b) $N = 5$, $\alpha = 0.05$, $t_s = 0.50\tau_2$. (c) $N = 6$ $(d = 64)$, $\alpha = 0.55$, $t_s = 0.08\tau_2$. (d) $N = 6$, $\alpha = 0.05$, $t_s = 0.50\tau_2$. The vertical gray lines mark t_s at which the reset channel is removed. Inset: schematic representation of the dissipative TFIM, with each spin coupled to a thermal bath.

of the Lindbladian. A larger α implies that ρ_0 is farther from ρ_{ss}. Here, ρ_{ss} is nontrivial, with nonzero entanglement. Moreover, as shown in [59], for a large fraction of randomly chosen pure initial states our protocol with $\rho_\delta = \mathbb{I}/d$ still yields substantial acceleration, further demonstrating its robustness. The parameters are $t_s = 0.5\tau_2$ and $t_s = 0.08\tau_2$ with $\tau_2 = 1/|\mathrm{Re}\lambda_2|$, the relaxation timescale, total duration $T = 6\tau_2$, $\alpha = 0.55/0.05$. Notably, t_s is chosen arbitrarily rather than being determined *a priori*. The accelerated relaxation processes under different reset rates are shown in Fig. 2. We also examine cases with different parameters in [59], where acceleration is consistently observed. In Figs. 2(b) and 2(d), genuine Mpemba effects are observed, where the system is first driven away from the steady state by the reset protocol compared to the no-reset case and then reaches stationarity faster. The existence of such an effect broadens the range of usable reset states.

We further verify in [59] that the same protocol significantly accelerates relaxation in another many-body setting, the Dicke model, again using $\rho_\delta = \mathbb{I}/d$.

Some remarks are in order regarding experimental implications. In practice, the precise optimal duration t_s is generally unknown in advance, but as emphasized before, fine-tuning is unnecessary. The TFIM example directly supports this conclusion: any choice of t_s shorter than the intrinsic relaxation time already leads to a clear acceleration. Additionally, the reset channel in the example corresponds to a depolarizing channel, which is conceptually straightforward to implement, e.g., by applying an isotropic white-noise field or coupling an infinite-temperature reservoir.

Experimental feasibility in general cases—The TFIM example illustrates the power of our protocol: the acceleration of relaxation via a depolarizing channel has the potential to be applied to complex quantum platforms. However, in general, it may be necessary to choose other ρ_δ to achieve acceleration. For few-body systems (such as single Rydberg atoms or other multilevel emitters), the protocol may be implemented with quantum reservoir engineering. Such techniques have become increasingly routine for tailoring dissipative dynamics in diverse quantum systems, including atoms [73–75], trapped ions [47,76], superconducting circuits [49,73,77], and optomechanical setups [78–80]. For quantum many-body systems, an exact implementation via reservoir engineering generally requires high-order interactions. Such engineered N-body dissipators have been put forward theoretically [81,82] and realized experimentally up to four-body interactions [83].

A complementary route is to approximate the reset channel: combine any available state-preparation protocol for ρ_δ, whether reservoir engineering, measurement-based feedback, or other methods, with Trotterization [60,61,84] (see End Matter). This strategy sidesteps the need to engineer reset jump operators directly and therefore avoids high-order couplings even in many-body systems. The main experimental challenge then reduces to preparing ρ_δ. Encouragingly, fast and high-fidelity state preparation has been reported across diverse platforms [62,75,85,86]. Notably, Ref. [62] realized state preparation in a 35-spin TFIM with Trotterization. This suggests that our protocol may be testable and applicable on existing platforms.

To connect our theoretical protocol more directly with experiments, we provide in [59] proposals that relate r to experimentally tunable parameters, offering realistic paths toward implementation.

Concluding remarks—We have introduced a general framework for accelerating relaxation in open quantum systems via temporary reset, applicable to arbitrary initial states. The proposed protocol can suppress multiple relaxation modes simultaneously, enabling enhanced control over the relaxation dynamics. As a practical example, we demonstrated that the relaxation of a TFIM can be significantly accelerated by using a simple depolarizing channel. This example highlights a powerful feature of our protocol: leveraging easily prepared states to accelerate the preparation of states that are difficult to reach. Furthermore, introducing temporary dephasing noise—a different type of channel—can also accelerate relaxation (End Matter). This suggests the broader applicability of our central idea: temporarily coupling the system to various quantum channels may provide a general route to enhanced relaxation. In future studies, our approach could be extended to certain non-Markovian and Floquet dissipative systems (e.g., using ideas in [9,87]). Overall, our results establish temporary reset as a powerful and experimentally feasible tool for controlling relaxation timescales in open quantum dynamics.

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Data availability—The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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End Matter

Derivation of the modified Lindbladian $\mathcal{L}_r$—Here, we derive Eq. (4) of the main text. Let $\{|\phi_i\rangle\}_{i=1}^{d}$ be an orthonormal basis, satisfying $\sum_i |\phi_i\rangle\langle\phi_i| = \mathbb{I}$. The reset state is generally written as $\rho_\delta = \sum_\alpha p_\alpha |\psi_\alpha\rangle\langle\psi_\alpha|$. Introduce jump operators

$$J^r_{i,\alpha} = \sqrt{r p_\alpha}|\psi_\alpha\rangle\langle\phi_i|, \quad i = 1, \ldots, d.$$

The associated Lindblad term is

$$\mathcal{D}_r(\rho) = \sum_{i,\alpha}\left(J^r_{i,\alpha}\rho J^{r\dagger}_{i,\alpha} - \frac{1}{2}\{J^{r\dagger}_{i,\alpha}J^r_{i,\alpha}, \rho\}\right). \tag{A1}$$

Direct evaluation gives

$$\sum_{i,\alpha} J^r_{i,\alpha}\rho J^{r\dagger}_{i,\alpha} = r\rho_\delta \mathrm{Tr}[\rho]$$

$$-\frac{1}{2}\sum_{i,\alpha}\{J^{r\dagger}_{i,\alpha}J^r_{i,\alpha}, \rho\} = -r\rho,$$

where $\mathrm{Tr}[\rho] = \sum_i \langle\phi_i|\rho|\phi_i\rangle$, $\sum_i |\phi_i\rangle\langle\phi_i| = \mathbb{I}$ and $\langle\psi_\alpha|\psi_\alpha\rangle = 1$ have been used. This yields $\mathcal{D}_r(\rho) = r[\mathrm{Tr}[\rho]\rho_\delta - \rho]$. Hence, we arrive at Eq. (4):

$$\mathcal{L}_r(\rho) = \mathcal{L}(\rho) + r\mathrm{Tr}[\rho]\rho_\delta - r\rho. \tag{A2}$$

Derivation of the modified coefficients—Here, we derive Eq. (7), the explicit form of the modified coefficients with reset. Acting the modified semigroup $e^{t\mathcal{L}_r}$ on both sides of

$$\rho_0 = \rho_{\mathrm{ss}} + \sum_{k=2}^{d^2} \mathrm{Tr}(L_k^\dagger \rho_0) R_k \tag{B1}$$

yields

$$\rho^r(t) = e^{t\mathcal{L}_r}\rho_0 = e^{t\mathcal{L}_r}\rho_{\mathrm{ss}} + \sum_{k=2}^{d^2} \mathrm{Tr}(L_k^\dagger \rho_0) e^{(\lambda_k - r)t} R_k. \tag{B2}$$

To proceed, we use an explicit form of the steady state under reset dynamics [44]

$$\rho_{\mathrm{ss}}^r = \lim_{t\to\infty}\rho^r(t) = \rho_{\mathrm{ss}} + \sum_{k=2}^{d^2}\frac{rd_k}{r-\lambda_k}R_k. \tag{B3}$$

Then, expressing the steady-state condition $e^{t\mathcal{L}_r}\rho_{\mathrm{ss}}^r = \rho_{\mathrm{ss}}^r$ with Eq. (B3) and applying Eq. (5), we get

$$e^{t\mathcal{L}_r}\rho_{\mathrm{ss}} + \sum_{k=2}^{d^2}\frac{rd_k}{r-\lambda_k}e^{(\lambda_k-r)t}R_k = \rho_{\mathrm{ss}} + \sum_{k=2}^{d^2}\frac{rd_k}{r-\lambda_k}R_k,$$

from which we solve $e^{t\mathcal{L}_r}\rho_{\mathrm{ss}}$. Substituting $e^{t\mathcal{L}_r}\rho_{\mathrm{ss}}$ into Eq. (B2) and comparing with Eq. (6), we identify modified coefficients $c_k^r(t)$ defined in Eq. (7).

Realizing the reset protocol with arbitrary state preparation processes—We now show that any state preparation protocol could approximately realize the reset channel in our scheme via a Trotterization-based construction. Notably, Trotterization has already been implemented experimentally in superconducting quantum circuits [88], IBM Quantum's hardware [84], and even quantum many-body platforms [62] to generate dissipative dynamics, such as the dissipative TFIM. Moreover, tensor network methods offer promising avenues for extending Trotterization to more complex open quantum systems [60]. We first consider the case where the original dynamics is absent $[\mathcal{L}(\rho)=0]$ to gain insight. In this case, the generator corresponding to the reset channel simply reads as

$$\dot{\rho}(t) = \mathcal{R}[\rho(t)] = r[\rho_\delta - \rho(t)]. \tag{C1}$$

Thus, a direct integration shows that applying the reset channel for a time t_s maps an initial state ρ_0 to a final state $\rho(t_s)$ according to

$$\rho(t_s) = e^{\mathcal{R}t_s}[\rho_0] = p\rho_0 + (1-p)\rho_\delta, \quad p \coloneqq e^{-rt_s}, \tag{C2}$$

that is, a reset channel with rate $r=-\ln p/t_s$ is effectively equivalent to a state preparation protocol for ρ_δ, with success probability $1-p$. p is a classical probability generated beforehand. One simply flips a biased coin (or, equivalently, performs any local measurement that yields the desired classical randomness): with probability $1-p$ the system is driven into ρ_δ by any suitable operations and with the probability p it is left untouched. Since there are no restrictions on the preparation method, once ρ_δ is locally preparable (e.g., a separable state), this operation can be implemented purely with local operations, even for many-body systems. Specifically, when $\rho_\delta = \mathbb{I}/d$, Eq. (C2) is a depolarizing channel, which can be directly realized by a unitary 2-design [59,63] without flipping a coin.

When $\mathcal{L}\neq 0$, we apply the Lie-Trotter formula [89]

$$e^{(\mathcal{L}+\mathcal{R})t} = \lim_{n\to\infty}\left(e^{\mathcal{R}\frac{t}{n}}e^{\mathcal{L}\frac{t}{n}}\right)^n \tag{C3}$$

to realize the protocol. With this formula, the reset protocol during $[0,t_s]$ can be approximated as

$$e^{(\mathcal{L}+\mathcal{R})t_s}[\rho_0] \approx (e^{\mathcal{R}\delta t}e^{\mathcal{L}\delta t})^n[\rho_0], \tag{C4}$$

where $\delta t \coloneqq t_s/n \ll 1$. Using Eq. (C2), we have

$$e^{\mathcal{R}\delta t}[\rho] \approx (1-p_s)\rho + p_s\rho_\delta, \tag{C5}$$

which can be interpreted as a state preparation mapping from ρ to ρ_δ, with success probability $p_s \coloneqq r\delta t$. Practically, r is determined by p_s and δt, whose values are set by experimentalists.

Thus, the reset protocol may be achieved experimentally by performing the aforementioned state preparation operations stroboscopically: at each discrete time point $t_i = i\delta t, (i=1,\ldots,n)$ within $[0,t_s]$, one applies the state preparation mapping from $\rho(t_i)$ to ρ_δ with a given probability p_s. There are no constraints on the preparation time τ_{prep} for ρ_δ, but on average it adds a cost of $np_s\tau_{\mathrm{prep}}$ to the tailored relaxation timescale. Therefore, it is desirable to choose a ρ_δ that can be prepared efficiently, e.g., the maximally mixed state in our TFIM example. This state is convenient to prepare: coupling to an infinite-temperature bath can achieve it in $\mathcal{O}(\ln N)$ time [90] with minimal resources in a N-body system—typically negligible compared to the system's intrinsic relaxation time. Local control methods could in principle realize it in constant time, though with resources scaling with N.

This Trotterization-based method admits a natural interpretation at the level of stochastic trajectories: for a Poissonian reset process with rate r, the system is stochastically reset to the target state with probability $r\delta t$, in each small interval δt.

The approximation error is of order $\mathcal{O}(t_s^2/n)$ [61,64,65]. Explicitly, we establish a rigorous upper bound [59]:

$$||e^{(\mathcal{L}+\mathcal{R})t_s} - (e^{\mathcal{R}\delta t}e^{\mathcal{L}\delta t})^n|| \le \frac{t_s^2}{2n}|||[\mathcal{L},\mathcal{R}]|||, \tag{C6}$$

where $||\cdot||$ denotes any norm that is contractive under Lindbladian evolution, such as the trace norm or the diamond norm. This bound holds for arbitrary Lindbladian $\mathcal{L}$ and $\mathcal{R}$. In our specific case,

$$[\mathcal{L},\mathcal{R}]\rho \coloneqq [\mathcal{L}\mathcal{R}-\mathcal{R}\mathcal{L}]\rho = r\mathrm{Tr}(\rho)\mathcal{L}(\rho_\delta), \tag{C7}$$

where $\mathrm{Tr}[\mathcal{L}(\rho)]=0$ is used. Hence, a large n is not required when $t_s \ll 1$, $rt_s \ll 1$ or ρ_δ is close to ρ_{ss}. A small t_s with moderate r offers a practical regime that may yield

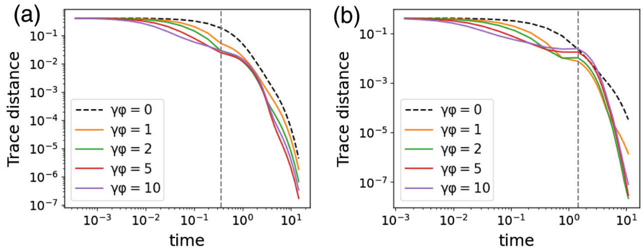

FIG. 3. Dephasing noise induces acceleration of relaxation in a 5-site TFIM. Parameters: (a) $J = 1.0, g = 2.0, \gamma = 0.5, \beta = 0.1J$, $t_s = 0.2\tau_2$ (b) $J = 1.0, g = 2.0, \gamma = 0.5, \beta = 0.1J, t_s = 0.8\tau_2$.

substantial acceleration and simultaneously reduce the required number of Trotter steps.

Additionally, higher-order approximations could be used to reduce the error for fixed n. For instance, the second-order Suzuki-Trotter formula $e^{(\mathcal{L}+\mathcal{R})t_s} \approx (e^{\mathcal{L}\delta t/2} e^{\mathcal{R}\delta t} e^{\mathcal{L}\delta t/2})^n$ suppresses the error to $\mathcal{O}(t_s^3/n^2)$ [61,64].

Dephasing noise can accelerate the relaxation of the TFIM—Here, we add dephasing noise along a single axis (typically the z axis) to the TFIM. The corresponding jump operators are $L_i^{(\phi)} = \sqrt{\gamma_\phi}\sigma_i^z, i = 1, \ldots, N$. This contributes to the total Lindbladian via an extra term

$$\mathcal{L}_\phi[\rho] = \gamma_\phi \sum_{i=1}^{N} (\sigma_i^z \rho \sigma_i^z - \rho), \tag{D1}$$

which could be interpreted as a partial and local reset channel (only the z axis is affected, and only local one-body jump operators are involved). As shown in Fig. 3, the relaxation is accelerated significantly.