

Blueprint for experiments exploring the quantum recurrence theorem on a coupled multiqubit system

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The quantum form of the Poincaré recurrence theorem stipulates that a system with a time-independent Hamiltonian and discrete energy levels returns arbitrarily close to its initial state in a finite time. Qubit systems, being highly isolated from their dissipative surroundings, provide a possible experimental test bed for studying this theoretical construct. Here, we investigate an N -qubit system, weakly coupled to its environment. We present quantitative analytical and numerical results on both the revival probability and time, and demonstrate that the system indeed returns arbitrarily close to its initial state, but in a time exponential in the number of qubits N , with revival times that are astronomically large for systems with just a few tens of qubits. Given the lifetimes achievable in present-day superconducting multiqubit systems, we propose a realistic experimental test of this theory and its size scaling. This provides insights into how thermalization emerges in isolated quantum systems.

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Introduction. Poincaré recurrence, both in its classical [1–4] and quantum [5–10] forms, is relevant to discussions of thermalization, quantum chaos [11], and the foundations of quantum thermodynamics [12]. It places strict constraints on the long-time behavior of isolated quantum systems, in spite of their apparent equilibration over shorter timescales [13–15]. Classical recurrence theory was explored in the 19th century, whereas the quantum recurrence theory (QRT) is known only since the 1950s, from when it has been applied to time-independent quantum mechanical systems with discrete energy eigenstates [16,17]. To our knowledge, there are no experimental demonstrations of the quantum recurrence theorem or its time scaling with system size.

Here, we investigate QRT in a coupled many-body quantum system. We consider first a quantum two-level system, a qubit, that is coupled to a quantum mechanical environment with a finite number N of degrees of freedom. A natural question is how the Poincaré recurrence time for the qubit scales with the size of this environment for a given return fidelity. This relates intrinsically to the problems of thermalization in isolated quantum systems [18–20] and Loschmidt echoes in quantum many-body dynamics [21], as well as to questions

of what forms a heat bath in a quantum system [22]. These issues are thought to be difficult to address in an experimental quantum system, due to the confounding dissipative effects from the environment. A theoretical analysis of the Poincaré recurrence in this system provides insights into the small N behavior, which we use to then explore Poincaré recurrence in an experimentally realizable construct of $N \sim \mathcal{O}(10)$ superconducting qubits, in which signatures of Poincaré recurrence should be observable over reasonable timescales, even in the presence of realistic levels of environmentally induced dissipation.

In a basic configuration, a central “test” qubit is coupled to an ensemble of N “environmental” qubits, as shown in Fig. 1(a), where we vary the energy splittings and coupling strengths of the qubits. We find that, as expected, the Poincaré recurrence time for the test qubit scales exponentially with N for a given range of splittings and coupling strengths. We derive analytical results for this system in the weak linear coupling regime and perform numerical analyses that are not restricted to the weak coupling condition. With just a few tens of environmental qubits, the system no longer recovers to its initial state on any realistic timescale.

We then build a multiscale model, where the isolated system of $N + 1$ qubits is connected to a much more complex environmental bath responsible for the usual decoherence of the qubits, with couplings that yield realistic qubit relaxation rates. This allows analysis relevant to an actual experimental system and provides a good measure as to what extent and over what timescale the $N + 1$ model qubits can be treated as an isolated quantum system. Finally, we relate our results to recently published experimental platforms [23,24] and find that, given current relaxation and decoherence rates,

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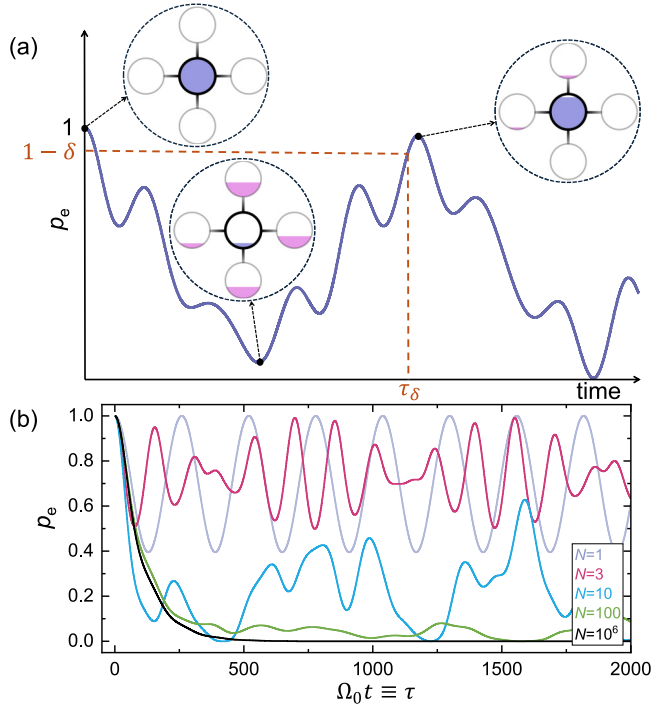


FIG. 1. Model for Poincaré revivals. (a) At time $t = 0$, a central “test” qubit is initialized in its excited state, while it is coupled to N “environmental” qubits (here $N = 4$), all of which are initially in their ground states. The solid line illustrates the resulting time evolution of the excited-state population p_e of the central qubit. The accompanying schematic provides a conceptual visualization of the qubit configurations at various points in time. Here, τ_δ is the first passage time for the test qubit population to return to within a threshold $\delta = 1 - p_e(\tau_\delta)$ of its initial state. Poincaré revival time τ_P is the expectation value of τ_δ averaged over different configurations of system parameters. (b) Example time traces of the excited-state population p_e of the central qubit for different numbers N of environmental qubits. For $N = 1$, we see the expected sinusoidal Rabi oscillations, but with increasing N the oscillations become less regular and weaker until for $N = 10^6$, we see only exponential decay.

experimental tests are fully feasible for multiqubit systems with tunable couplings and energies [25–27].

The main theoretical results can be obtained from the analysis of an archetypal synthetic system described by the standard Hamiltonian, $\hat{H} = \hat{H}_S + \hat{H}_E + \hat{H}_c$, where $\hat{H}_S = \hbar\Omega_0\hat{a}^\dagger\hat{a}$ represents the Hamiltonian of the test qubit with energy splitting $\hbar\Omega_0$, $\hat{H}_E = \sum_{j=1}^N \hbar\Omega_j\hat{b}_j^\dagger\hat{b}_j$ denotes the Hamiltonian of the N environmental qubits with energies $\hbar\Omega_j$, and $\hat{H}_c = \sum_{j=1}^N g_{0j}(\hat{a}^\dagger\hat{b}_j + \hat{b}_j^\dagger\hat{a})$ models the system-environment interaction, assumed to be linear with a coupling strength g_{0j} between the test qubit and the j th environmental qubit; $\hat{a}$ and $\hat{b}_j$ are the annihilation operators for the corresponding qubits, respectively. In a real experimental system, both the test and environmental qubits will be coupled to a complex and uncontrolled bath, to be considered later in this paper.

We start by solving the Schrödinger equation, $i\hbar\partial_t|\psi_I(t)\rangle = \hat{H}_{c,I}(t)|\psi_I(t)\rangle$, for the system in the single excitation subspace, including the test qubit and the N environmental qubits. Here, the subscript I indicates that the coupling Hamiltonian and the wave function are expressed in

the interaction picture. We then obtain a set of time-evolution equations [28] for which an analytical solution can be found by proceeding iteratively. Figure 1(b) presents numerical examples of how the excited-state population p_e of the test qubit evolves when it is in contact with N environmental qubits. For small N , we observe revivals, with individual traces demonstrating an almost periodic function [29,30]. For large N , these revivals disappear entirely. In this figure, $\Delta\Omega$ is the energy spread of the environmental qubits, with $\Delta\Omega/\Omega_0 = 0.1$, and we parametrize the coupling by the rate Γ_0 related to g_{0j} , here normalized by $\hbar\Omega_0$, through $\langle g_{0j}^2 \rangle = \Delta\Omega\Gamma_0/(2\pi\Omega_0^2N)$ [28], with $\Gamma_0 = 0.01\Omega_0$, and the coupling strengths g_{0j} are distributed uniformly from 0 up to a maximum value determined by Γ_0 .

Quantum recurrence. Starting with the isolated system with N environmental qubits, as illustrated in Fig. 1(a) for $N = 4$, we prepare the test qubit in its excited state and monitor its evolution as it interacts with the environmental qubits, which are initialized in their ground states.

We first demonstrate that monitoring the excited-state population of the test qubit serves as a measure of quantum revivals. We define the basis formed by the single-excitation subspace $|0\rangle \equiv |100\dots 0\rangle$ and $|j\rangle \equiv |000\dots 1(j\text{th})\dots 0\rangle$, where the first entry corresponds to the test qubit followed by the entries corresponding to each of the N environmental qubits. Assuming the arrangement shown schematically in Fig. 1, at time $t = 0$ the state of the system is $|\psi(0)\rangle = |100\dots 0\rangle \equiv |0\rangle$.

After a time t , the state is given by

$$|\psi(t)\rangle = \mathcal{C}_0(t)|0\rangle + \sum_{j=1}^N \mathcal{C}_j(t)|j\rangle, \quad (1)$$

where $|\mathcal{C}_0(t)|^2 = p_e(t)$ and $|\mathcal{C}_j(t)|^2$ represent the population of the test qubit and the j th environmental qubit, respectively. We define the distance d of the state at t from the initial state as $d = \sqrt{\langle \delta\psi(t) | \delta\psi(t) \rangle}$, where $|\delta\psi(t)\rangle \equiv |\psi(t)\rangle - |\psi(0)\rangle$. We obtain $d^2 = 2(1 - |\mathcal{C}_0(t)|^2)/(1 + |\mathcal{C}_0(t)|^2) \leq 2(1 - p_e(t))$. Thus if $1 - p_e(t)$ is small, then d^2 is also small, justifying the use of $1 - p_e(t)$ as a measure of the recurrence with threshold δ as shown in Fig. 1(a). In this panel, we show the evolution for $p_e(t)$ for a single instance of $N = 4$ environmental qubits, and in panel (b) example time traces for a range of environment sizes N , evolving from Rabi-like oscillations for small N to exponential decay for large N , each trace for fixed values of the couplings and energies.

Analytic derivations. We next present analytical expressions for the central quantities governing the Poincaré revivals. We first demonstrate the exponential increase of the revival time τ_P in the large N limit for the system described above, averaging over random distributions of environmental qubit energies and couplings. Due to the central limit theorem, the probability distribution $p_N(\Delta)$ for $\Delta \equiv 1 - p_e(\tau)$ will have a Gaussian distribution $p_N(\Delta) = \frac{1}{\sqrt{2\pi s^2}} e^{-(\Delta - \langle \Delta \rangle)^2/(2s^2)}$, where $s^2 = \langle \Delta^2 \rangle - \langle \Delta \rangle^2$ is the variance. For N qubits with independent randomly distributed energies, we have $\langle \Delta \rangle = N\langle \Delta_j \rangle$ and $s^2 = N s_j^2$, where $\langle \Delta_j \rangle$ and s_j^2 are the mean and the variance for individual qubits, respectively. For small threshold

δ (i.e., near a recovery), we then have

$$p_N(\Delta) \propto e^{-\langle \Delta \rangle^2 / (2s^2)} \equiv e^{-bN}, \quad (2)$$

where $b = \langle \Delta \rangle^2 / (2s^2)$. The total probability $\mu_N(\delta)$ to stay below δ is then also $\propto e^{-bN}$. The Poincaré recovery time τ_P , i.e., the expectation value of the revival time τ_δ , on the other hand, is $\tau_P \propto \mu_N^{-1} = e^{bN}$, as will be discussed below.

For a more rigorous analytic derivation of μ_N and τ_P (see details in the Supplemental Material [31]), we focus on the dynamics of $\mathcal{C}_0(t)$, applying the given initial conditions and retaining only the first-order terms, yielding [28]

$$p_e(t) \equiv |\mathcal{C}_0(t)|^2 = 1 - 4 \sum_{j=1}^N \mathcal{G}_j^2 \frac{\sin^2[(\lambda_j - 1)\tau/2]}{(\lambda_j - 1)^2}, \quad (3)$$

where $\mathcal{G}_j \equiv g_{0j}/\hbar\Omega_0$, $\lambda_j \equiv \Omega_j/\Omega_0$, and $\tau = \Omega_0 t$. We may identify $\Delta_j = 4\mathcal{G}_j^2 \sin^2(\phi_j/2)/\eta_j^2$; thus, $p_e(\tau) \equiv 1 - \Delta$, with $\Delta = \sum_{j=1}^N \Delta_j$. Here, $\eta_j \equiv \lambda_j - 1$ and $\eta_j \tau \equiv \phi_j$. For recoveries $p_e(\tau) \sim 1$, all the (positive) terms Δ_j in the sum are very small. The weak coupling approximation is well justified if $\mathcal{G}_j \lesssim 10^{-3}$ (see Supplemental Material for an illustration [31]). At sufficiently long times τ , ϕ_j is randomly and uniformly distributed in the interval $[-\pi, \pi]$ modulo 2π . The randomness of phase holds for times $\tau \gg \sqrt{\delta/\mathcal{G}}$ [32]. For simplicity, we assume that the energies η_j of the environmental qubits are uniformly distributed around that of the test qubit, in the interval $-\Delta\Omega/(2\Omega_0) < \eta_j < \Delta\Omega/(2\Omega_0)$. We then obtain the distribution of Δ_j for the case of equal coupling to all qubits $\mathcal{G}_j \equiv \mathcal{G}$, by calculating the probability via $P(\Delta_j) = \int d\phi_j \int d\eta_j p(\phi_j) p(\eta_j)$. Here, $p(\phi_j)$ and $p(\eta_j)$ are the uniform probability density functions. In the regime $\Delta_j < \Delta_m = 16\mathcal{G}^2/(\Delta\Omega/\Omega_0)^2$, we have

$$P(\Delta_j) = \frac{2}{\pi} \left[\sqrt{\frac{\Delta_m}{\Delta_j}} - 1 - \sqrt{\frac{\Delta_m}{\Delta_j}} + \arcsin\left(\sqrt{\frac{\Delta_j}{\Delta_m}}\right) \right]. \quad (4)$$

The probability density is given by $p(\Delta_j) = dP(\Delta_j)/d\Delta_j$. For $\Delta_j \ll \Delta_m$, we have

$$p(\Delta_j) \simeq \frac{1}{\mathcal{T}} \frac{1}{\sqrt{\Delta_j}}, \quad (5)$$

with $\mathcal{T} = 8\pi\mathcal{G}/(\Delta\Omega/\Omega_0)$. Equation (5) allows us then to calculate the corresponding probability density for Δ as

$$p_N(\Delta) = \frac{1}{(64\pi)^{\frac{N}{2}} \Gamma(\frac{N}{2})} \left(\frac{\Delta\Omega/\Omega_0}{\mathcal{G}} \right)^N \Delta^{\frac{N}{2}-1}. \quad (6)$$

Then, the revival probability $\mu_N(\delta) = \int_0^\delta p_N(\Delta) d\Delta$ for being within the threshold, that is, in the range $1 - p_e(\tau) < \delta$, is

$$\mu_N(\delta) = \frac{1}{(64\pi)^{\frac{N}{2}} \Gamma(\frac{N}{2} + 1)} \left(\frac{\Delta\Omega/\Omega_0}{\mathcal{G}} \right)^N \delta^{\frac{N}{2}}. \quad (7)$$

The distributions $p_N(\Delta)$ and $\mu_N(\delta)$ for a few values of N are shown in Fig. 2 together with the Gaussian approximation. This demonstrates the validity range of Eqs. (6) and (7), limited by the condition $\Delta_j \ll \Delta_m$.

Next, we compare the analytical expression, Eq. (7), with the full Schrödinger equation solution and the outcome of

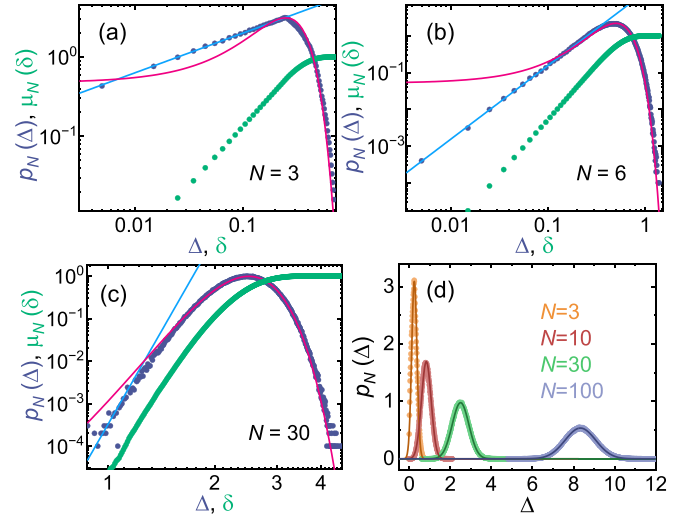


FIG. 2. Probability density distributions $p_N(\Delta)$ of $\Delta \equiv \sum_{i=1}^N \Delta_i$ for (a) $N = 3$, (b) $N = 6$, and (c) $N = 30$. For illustration, we set $\mathcal{T} = 1$ here. The magenta line is the Gaussian distribution, the dark blue symbols show the numerical stochastic values (histogram), and the aqua line is from Eq. (6), valid for small Δ . The green symbols show the numerically calculated cumulative histogram $\mu_N(\delta) = \int_0^\delta p_N(\Delta) d\Delta$. Panel (d) shows the same distribution on a linear scale for four different values of N .

Eq. (3), both averaged over a long time interval. These results are shown in Fig. 3(a) for small N for the parameters $\mathcal{G} = 0.001$, $\Delta\Omega/\Omega_0 = 0.1$, and $\delta = 0.001$, within the validity range of the analytical expression. We see that these results coincide in a satisfactory way and μ_N indeed decays exponentially, even for these small values of N .

To evaluate the actual revival time τ_P , we find that the rate r to cross the threshold at $p_e = 1 - \delta$ is given by

$$r = p_N(\delta) \langle R \rangle_+, \quad (8)$$

where $\langle R \rangle_+ \equiv \int_0^\infty dR p_R(R) R$ is the expectation value of $R \equiv -d\Delta/d\tau$ for positive slopes at $\Delta = \delta$ and $p_R(R)$ is the probability density of R . Incorporating the survival probability $P_s(\tau)$, we have $dP_s(\tau)/d\tau = -rP_s(\tau)$, which yields the revival time $\tau_P = r^{-1}$ as

$$\tau_P = \frac{\sqrt{2\pi} \Gamma(\frac{N+1}{2})}{N^{3/2} \Gamma(\frac{N}{2})} \frac{\sqrt{\delta}}{\mathcal{G}} \mu_N(\delta)^{-1}. \quad (9)$$

Panels (c) and (d) in Fig. 3 display the results for the Poincaré revival time τ_P as a function of the environment size N . In Fig. 3(c), we show the comparison of the numerical and analytical results, using the same parameters as in Fig. 3(a) for $\mu_N(\delta)$. We see an almost exponential increase of $\tau_P \propto \mu_N^{-1}$ with N , even for these small values of N , with good agreement between numerical and analytical results. The numerical results, outside the weak-coupling regime, shown in Fig. 3(d), further illustrate the exponential dependence on N .

An experimentally realistic model. A superconducting multiqubit platform that can be adapted to a test of the Poincaré revival [23] is shown schematically in Fig. 4(a). We now include in our model the expected environmental decoherence, by introducing a bath of two-level systems (TLSs) interacting

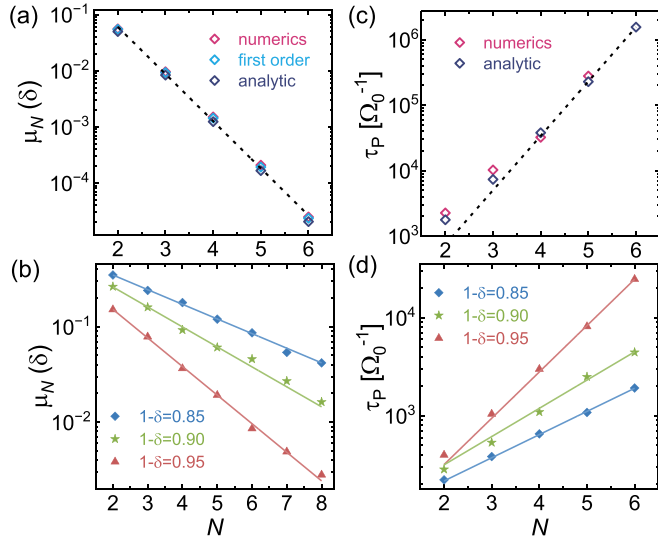


FIG. 3. Revival probability $\mu_N(\delta)$ to have $p_e(\tau) > 1 - \delta$ in the long-time limit, and the Poincaré time τ_P . (a) Comparison between the full numerical solution of the Schrödinger equation and the first-order approximation given by Eq. (3), as well as the analytical solution given by Eq. (7). The dashed line represents an exponential fit $\mu_N \propto e^{-1.93N}$. The parameters for this figure are $\Delta\Omega/\Omega_0 = 0.1$, $\mathcal{G} = 0.001$, and $\delta = 0.001$. (b) Numerical results for $\mu_N(\delta)$ for three values of δ . In each case, $\mu_N(\delta)$ exhibits an exponential decay with N , with the corresponding fits $e^{-0.35N}$, $e^{-0.48N}$, and $e^{-0.69N}$, ordered from top to bottom and shown by solid lines. The parameters used in this panel are $\Delta\Omega/\Omega_0 = 0.1$ and $\Gamma_0 = 0.01\Omega_0$, where $\mathcal{G}_j$ is uniformly distributed between 0 and its maximum value determined by Γ_0 . (c) Comparison of Poincaré revival time τ_P from the full numerical results and the analytical solution [Eq. (9)]. Dashed line represents an exponential fit of the form $\tau_P \propto \exp(bN)$ with $b = 1.91$. Parameters are as in panel (a). (d) Numerical results for the dependence of τ_P on N for different threshold values of δ . The data demonstrate the exponential increase predicted by the model. The parameters in this panel are as in panel (b).

with the $N + 1$ test and environmental qubits. The full Hamiltonian of the system in Fig. 4(b) is given by

$$\begin{aligned} \hat{\mathcal{H}} = & \hbar\Omega_0\hat{a}^\dagger\hat{a} + \sum_{i=1}^N \hbar\Omega_i\hat{b}_i^\dagger\hat{b}_i + \sum_{j=1}^M \hbar\omega_j\hat{c}_j^\dagger\hat{c}_j \\ & + \sum_{k=1}^N g_{0k}(\hat{a}^\dagger\hat{b}_k + \hat{b}_k^\dagger\hat{a}) + \sum_{l \neq m} g_{lm}\hat{b}_l^\dagger\hat{b}_m \\ & + \sum_{n=1}^M \gamma_{0n}(\hat{a}^\dagger\hat{c}_n + \hat{c}_n^\dagger\hat{a}) + \sum_{p,q} \gamma_{pq}(\hat{b}_p^\dagger\hat{c}_q + \hat{c}_q^\dagger\hat{b}_p). \end{aligned} \quad (10)$$

The first three terms describe the noninteracting Hamiltonian of the test qubit, the N environmental qubits, and a bath of M two-level systems with energies $\hbar\omega_j$. The fourth and fifth terms couple the test qubit and environmental qubits to one another. The final two terms couple the test and environmental qubits to the M TLSs, the latter with annihilation operators $\hat{c}_k$. In Eq. (10), g_{lm} denotes the coupling constant between the l th and m th environmental qubit; γ_{0n} the coupling constant between the test qubit and the n th TLS; and γ_{pq} the coupling between the p th environmental qubit and the q th bath TLS.

We use a Fock state basis $|0\rangle \equiv |1, 00 \dots 0, 00 \dots 0\rangle$, $|i\rangle \equiv |0, 00 \dots 1(i\text{th}) \dots 0, 00 \dots 0\rangle$, and $|\bar{j}\rangle \equiv |0, 00 \dots 0, 00 \dots 1(j\text{th}) \dots 0\rangle$, where the basis vectors are coded such that the first entry corresponds to the central qubit, the next N entries to each of the environmental qubits, and the third group, running from $2 + N$ to $2 + N + M$, to the M bath TLSs. For simplicity, we ignore the TLS-TLS coupling in the model. These three sets are separated by commas in the ket vector for clarity. Based on this construction, the state $|\psi(t)\rangle$ can be written as

$$|\psi(t)\rangle = \mathcal{C}_0|0\rangle + \sum_{i=1}^N \mathcal{C}_i|i\rangle + \sum_{j=1}^M D_j|\bar{j}\rangle. \quad (11)$$

In the interaction picture, the dynamics of the system are governed by the coupling terms via the perturbation $\hat{V}_I(t)$, which arises from the last four terms of the Hamiltonian (10). We again solve the Schrödinger equation, $i\hbar\partial_t|\psi_I(t)\rangle = \hat{V}_I(t)|\psi_I(t)\rangle$, for the complete system. This allows us to determine how the state of the entire system evolves over time (details given in the Supplemental Material [31]).

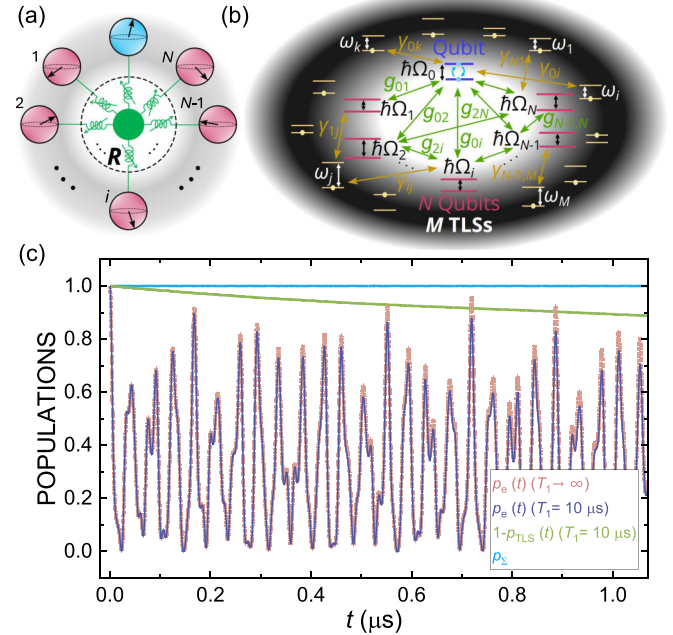


FIG. 4. A possible experimental setup consists of $N + 1$ qubits (one test qubit and N environmental qubits) coupled all-to-all via a superconducting router marked as R within the dashed circle [23]. (a) The energies of all the qubits $\hbar\Omega_j$ and their couplings g_{jk} are tunable. (b) A more realistic model of panel (a), taking into account the decoherence due to the external bath, which here is modeled by M bath TLSs. (c) Numerical results for system in panel (b). We use a model size $N = 5$, $M = 10000$, $T_1 = 10 \mu\text{s}$, with the parameters $\Delta\Omega/\Omega_0 = 0.02$ and $\Gamma_0/\Omega_0 = 4 \times 10^{-6}$. Purple and brown lines show the population $p_e(\tau) = |\mathcal{C}_0(\tau)|^2$ of the test qubit with and without relaxation, respectively. The green line is $\sum_{k=0}^N |\mathcal{C}_k(\tau)|^2$, showing the population in the $N + 1$ test and environmental qubits including relaxation, where decay results due to leakage to the TLS bath. Light blue line is $\sum_{k=0}^N |\mathcal{C}_k(\tau)|^2 + \sum_{j=1}^M |D_j(\tau)|^2$, which remains at unity, demonstrating conservation of probability in the closed system formed by all the qubits and TLSs.

To assess the experimental feasibility, we develop a numerical experiment with parameters based on an existing system [23]. Taking into account the finite relaxation time of the qubits (here using a conservative $T_1 \sim 10 \mu\text{s}$, where we have $T_1 = 1/\Gamma_0 = \frac{\Delta\Omega}{2\pi\Omega_B^2 M} \langle \gamma_{0n}^2 \rangle^{-1}$, with realistic couplings and qubit energy distributions, we calculate the $p_e(\tau)$ pattern shown in Fig. 4(c). For reference, we also plot results for a fully isolated system (with no coupling to the TLS bath) with identical qubit parameters. Experiments lasting a few microseconds should not be significantly affected by qubit decoherence, but for precise analysis, one likely needs to renormalize the threshold values based on leakage to the external TLS bath. This is indicated by the decreasing probability for the excitation to remain within the $N + 1$ qubit system, $\sum_{i=0}^N |\mathcal{C}_i(\tau)|^2$, as shown by the green line.

Conclusion. Whether and how a unitary quantum system thermalizes remains a fundamental open question, with direct relevance for scaling up quantum processors. The revival of an initial state in a many-body system, as in Poincaré recurrence, provides a fingerprint of the remnant quantum dynamics and is opposed to the concept of full thermalization. We have presented a quantitative analysis of the small N behavior of Poincaré recurrence in a quantum system, as well as a blueprint for investigating these quantum revivals experimentally. The proposed construct uses a collection of coupled qubits that form a quasi-isolated quantum system. Our analysis also includes the decohering effects of a realistic bath of TLSs, for which we derive explicit analytical and numerical results. Although our results are general, superconducting qubits are particularly suited for an experimental study, as these devices can be tailored in terms of their

energy spectra and couplings with an exceptional degree of control. Furthermore, present-day superconducting qubits are sufficiently isolated from the environment, reducing the bath decoherence to an acceptable level. Most importantly, individual qubits can be measured accurately in the time domain, allowing the direct experimental study of Poincaré recurrence.

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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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