

Programmable spectral symmetries in an anisotropic quantum Rabi simulator

Yu Liu

yuliu@iphy.ac.cn

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0009-0009-4223-1839>

Jia-Cheng Song

Beijing National Laboratory for Condensed Matter Physics, Institute of Physics, Chinese Academy of Sciences

Ming-Chuan Wang

Institute of Physics, Chinese Academy of Sciences

Ke-Xiong Yan

Fuzhou University

Yang He

Beijing National Laboratory for Condensed Matter Physics, Institute of Physics, Chinese Academy of Sciences

Yun-Hao Shi

Xi'an Jiaotong University

Wei-Ping Yuan

Institute of Physics, Chinese Academy of Sciences

Cheng-Lin Deng

Beijing Academy of Quantum Information Sciences

Li Li

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0009-0005-0785-6703>

Zhen-Ting Bao

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0009-0007-6924-4360>

Yutao Chen

Institute of Physics, Chinese Academy of Sciences

Xu-Yang Gu

Institute of Physics, Chinese Academy of Sciences

Tian-Ming Li

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0000-0002-8904-6317>

Gui-Han Liang

Institute of Physics, Chinese Academy of Sciences

Zheng-He Liu

Institute of Physics, Chinese Academy of Sciences

Wei-Guo Ma

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0009-0000-3639-2294>

Zhen-Yu Peng

Beijing National Laboratory for Condensed Matter Physics, Institute of Physics, Chinese Academy of Sciences

Shuai-Li Wang

Institute of Physics, Chinese Academy of Sciences

Yongxi Xiao

Institute of Physics, Chinese Academy of Sciences

Yi-Han Yu

Beijing National Laboratory for Condensed Matter Physics, Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0009-0000-8265-7409>

Jia-Chi Zhang

Institute of Physics, Chinese Academy of Sciences

Kui Zhao

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0009-0005-9908-8974>

Min-Xuan Zhou

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0009-0005-6012-2952>

Kaixuan Huang

Beijing Academy of Quantum Information Sciences <https://orcid.org/0000-0001-8845-2770>

Yu-Ran Zhang

School of Physics and Optoelectronics, South China University of Technology, Guangzhou 510640, China <https://orcid.org/0000-0002-6893-2053>

Yu-Xiang Zhang

Institute of Physics, Chinese Academy of Sciences

Zhongcheng Xiang

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0000-0002-0734-1217>

Dongning Zheng

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0000-0001-5895-0111>

Ye-Hong Chen

Fuzhou University

Kai Xu

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0000-0002-7776-2455>

Heng Fan

Institute of Physics, Chinese Academy of Sciences <https://orcid.org/0000-0002-5328-2449>

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Programmable spectral symmetries in an anisotropic quantum Rabi simulator

Jia-Cheng Song,^{1,2,*} Yu Liu,^{1,†} Ming-Chuan Wang,^{1,2,*} Ke-Xiong Yan,^{3,*} Yang He,^{1,2} Yun-Hao Shi,⁴ Wei-Ping Yuan,^{1,2} Cheng-Lin Deng,⁵ Li Li,^{1,2} Zhen-Ting Bao,^{1,2} Yutao Chen,^{1,2} Xu-Yang Gu,^{1,2} Tian-Ming Li,^{1,2} Gui-Han Liang,¹ Zheng-He Liu,^{1,2} Wei-Guo Ma,^{1,2} Zhen-Yu Peng,^{1,2} Shuai-Li Wang,^{1,2} Yong-Xi Xiao,^{1,2} Yi-Han Yu,^{1,2} Jia-Chi Zhang,^{1,2} Kui Zhao,⁵ Min-Xuan Zhou,^{1,2} Kaixuan Huang,⁵ Yu-Ran Zhang,⁶ Yu-Xiang Zhang,^{1,2,7} Zhongcheng Xiang,^{1,2,7,8} Dongning Zheng,^{1,2,7,8} Ye-Hong Chen,^{3,9,‡} Kai Xu,^{1,2,5,7,8,§} and Heng Fan^{1,2,5,7,8,¶}

¹*Beijing National Laboratory for Condensed Matter Physics,*

Institute of Physics, Chinese Academy of Sciences, Beijing 100190, China

²*School of Physical Sciences, University of Chinese Academy of Sciences, Beijing 100049, China*

³*Fujian Key Laboratory of Quantum Information and Quantum Optics, College of Physics and Information Engineering, Fuzhou University, Fuzhou, Fujian 350108, China*

⁴*School of Physics, Xi'an Jiaotong University, Xi'an 710049, China*

⁵*Beijing Key Laboratory of Fault-Tolerant Quantum Computing, Beijing Academy of Quantum Information Sciences, Beijing 100193, China*

⁶*School of Physics and Optoelectronics, South China University of Technology, Guangzhou 510640, China*

⁷*Hefei National Laboratory, Hefei 230088, China*

⁸*Beijing Key Laboratory of Advanced Quantum Technology, Beijing 100190, China*

⁹*Quantum Information Physics Theory Research Team,*

Center for Quantum Computing, RIKEN, Wako-shi, Saitama 351-0198, Japan

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The quantum Rabi model captures fundamental aspects of light-matter interaction, where symmetry dictates both spectra and dynamics. Over the past years, experiments have explored many of its nonperturbative properties, but have mostly focused on the isotropic limit, where rotating and counterrotating processes are locked together, leaving the broader symmetry landscape largely unexplored. Here we realize a programmable anisotropic quantum Rabi model in a superconducting processor, with independent control of the rotating and counterrotating couplings (g_1, g_2) and of a transverse bias ε . Continuous anisotropy tuning, combined with a duality mapping, gives access to the full parameter space from the Jaynes-Cummings to the anti-Jaynes-Cummings limits. In the deep-strong-coupling regime, we show that anisotropy reconstructs the spectrum and turns complete collapse-revival dynamics into incomplete revivals even near degeneracy. With adiabatic state preparation and joint tomography, we resolve an anisotropy-induced ground-state parity switch, a crossing that has no analogue in the isotropic model. We further observe selective tunnelling associated with hidden symmetry in biased Rabi models and track its anisotropic displacement within the same device. These results establish a controllable route to engineering nonperturbative light-matter Hamiltonians, where symmetry, spectrum, and dynamics can be programmed independently.

The quantum Rabi model (QRM) [1], a two-level system coupled to a single bosonic mode, is one of the simplest Hamiltonians, in which the light-matter interaction becomes genuinely quantum. It underlies phenomena ranging from coherent vacuum Rabi oscillations [2–4] to the generation of nonclassical states [5, 6]. When the interaction strength g is small compared with the mode frequency ω , the Jaynes-Cummings model (JCM) [7] provides an accurate description, because the total excitation number is conserved. This conservation law produces the characteristic $\sqrt{n}$ ladder [2, 8] and coherent matter-field exchange [5, 9, 10]. As the coupling approaches the mode frequency, the rotating-wave approximation (RWA) breaks down. The sys-

tem then enters a nonperturbative regime [3, 11–16], where excitation manifolds hybridize, vacuum dressing becomes pronounced [13, 17, 18], and intrinsically multi-frequency dynamics emerges [19, 20]. These effects define the ultrastrong-coupling (USC) regime, typically when $g/\omega > 0.1$, and the deep-strong-coupling (DSC) regime, often identified by $g/\omega \gtrsim 1$. Yet nonperturbative coupling strength is only one axis of the problem. The relative weight of rotating and counterrotating processes also determines the conserved quantities, selection rules, and level crossings of the model. Most experiments have nevertheless focused on the isotropic QRM [21–27], where this balance is fixed. As a result, central questions remain open: how does the anisotropy reorganize the spectrum, how do symmetry-protected crossings appear or disappear, and can hidden symmetries be detected dynamically within a single programmable device? Addressing these questions requires independent control of the two interaction channels, opening the access to a broader landscape beyond the RWA [28–32].

A natural framework for this control is the anisotropic quantum Rabi model (AQRM) [30], which extends the

* These authors contributed equally to this work.

† yuliu@iphy.ac.cn; These authors contributed equally to this work.

‡ yehong.chen@fzu.edu.cn

§ kaixu@iphy.ac.cn

¶ hfan@iphy.ac.cn

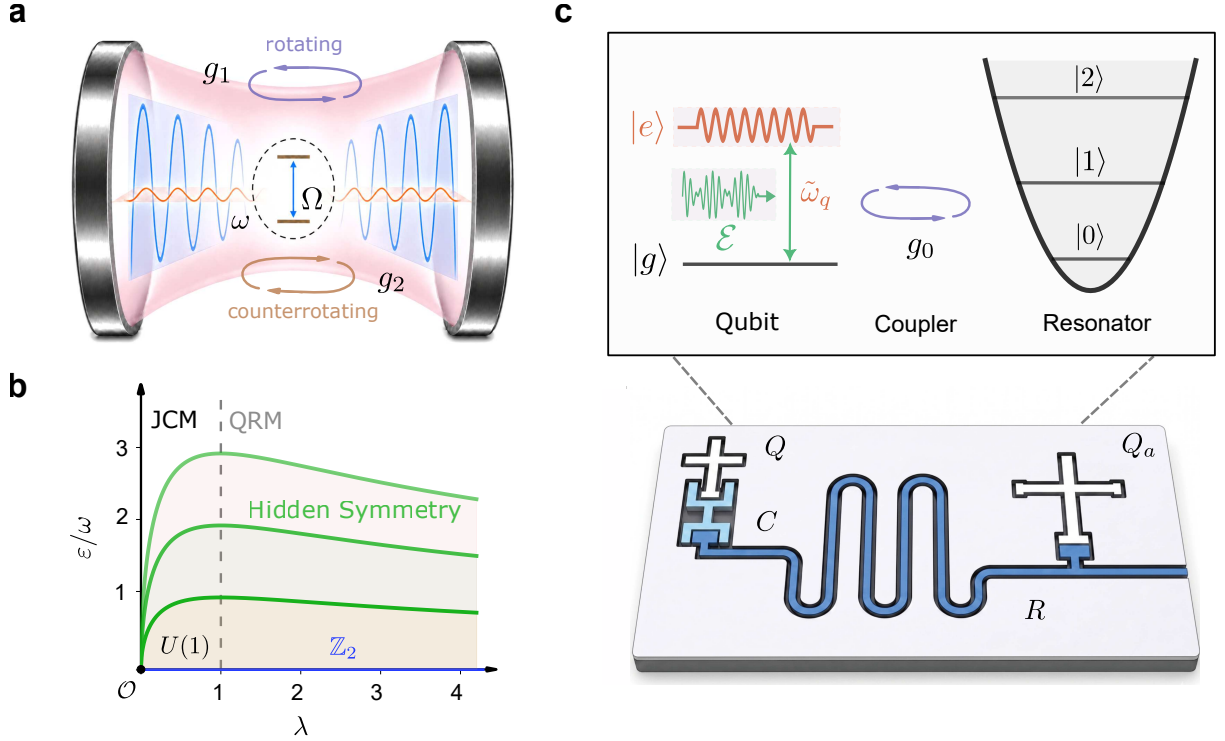


FIG. 1. **Device and theoretical model.** **a**, Schematic illustration of the anisotropic quantum Rabi model, describing a qubit coupled to two quadratures of the quantized light field. **b**, Symmetry structure of the AQRM in the parameter space $(\lambda, \varepsilon/\omega)$. The point $(0,0)$ corresponds to the $U(1)$ -symmetric JCM limit, the line $\varepsilon = 0$ preserves the $\mathbb{Z}_2$ parity symmetry. Even outside these symmetry manifolds, hidden symmetries can emerge at special parameter points. **c**, Schematic diagram of the experimental device and pulse sequence. The system qubit Q is coupled to the resonator while subject to longitudinal sinusoidal modulation and a multi-tone transverse drive. Together, these modulations realize an effective AQRM Hamiltonian.

QRM by coupling a two-level system to both quadratures of a bosonic mode, as illustrated in Fig. 1a. The AQRM provides a tunable setting for quantum criticality, quantum sensing, and nonclassical-state generation [31, 33–37]. The Hamiltonian takes the form (by setting $\hbar = 1$)

$$\hat{H} = \frac{\Omega}{2} \hat{\sigma}_z + \omega \hat{a}^\dagger \hat{a} + g_1 (\hat{H}_r + \lambda \hat{H}_{cr}) + \frac{\varepsilon}{2} \hat{\sigma}_x, \quad (1)$$

where $\hat{H}_r = \hat{a}^\dagger \hat{\sigma}_- + \hat{a} \hat{\sigma}_+$ and $\hat{H}_{cr} = \hat{a}^\dagger \hat{\sigma}_+ + \hat{a} \hat{\sigma}_-$ are the rotating and counterrotating interactions, respectively. Here, Ω , ω , and ε are the qubit frequency, mode frequency, and transverse bias, respectively. And $\lambda = g_2/g_1$ sets the relative strength of the two channels. Tuning λ interpolates between the JCM ($\lambda = 0$) and the isotropic QRM ($\lambda = 1$). The AQRM exhibits a layered symmetry structure. For $\lambda = 0$ and $\varepsilon = 0$, it has a continuous $U(1)$ symmetry associated with the conservation of the total excitation number $\hat{N} = \hat{a}^\dagger \hat{a} + (\hat{\sigma}_z + 1)/2$. For $\lambda \neq 0$ but $\varepsilon = 0$, this symmetry is reduced to a discrete $\mathbb{Z}_2$ parity symmetry generated by $\hat{P} = -(-1)^{\hat{a}^\dagger \hat{a}} \hat{\sigma}_z$ [38]. A finite bias generally breaks parity and turns protected crossings into avoided crossings, yet exact crossings can reemerge at the discrete commensurate biases associated with exceptional solutions and hidden symmetries [39–46]. This

symmetry structure is summarized in Fig. 1b. In addition to this rich structure and broad interest, direct experimental access to the AQRM and systematic probes of its spectra and dynamics have remained challenging, despite some theoretical discussions [47–49].

Here, we construct a programmable AQRM simulator based on the superconducting circuits, with independent control of the rotating and counterrotating channels (g_1, g_2) and of the bias ε . Together with a duality mapping, this tunability allows us to access the anisotropy continuously from $\lambda = 0$ to $\lambda = \infty$, thereby covering the full AQRM parameter space. Using the time-domain measurements, adiabatic state engineering and joint qubit-resonator tomography, we show how anisotropy and bias reshape the QRM in three connected ways. First, anisotropy modifies the level spacings in the DSC regime and converts complete collapse-revival dynamics into incomplete revivals. Second, it reorganizes the low-lying spectrum and produces a ground-state parity switch. Third, for the biased models, it shifts the dynamical signature of hidden symmetry in a way that follows the predicted anisotropic commensurability condition. Together these observations turn the AQRM from a theoretical interpolation into an experimentally programmable symmetry landscape.

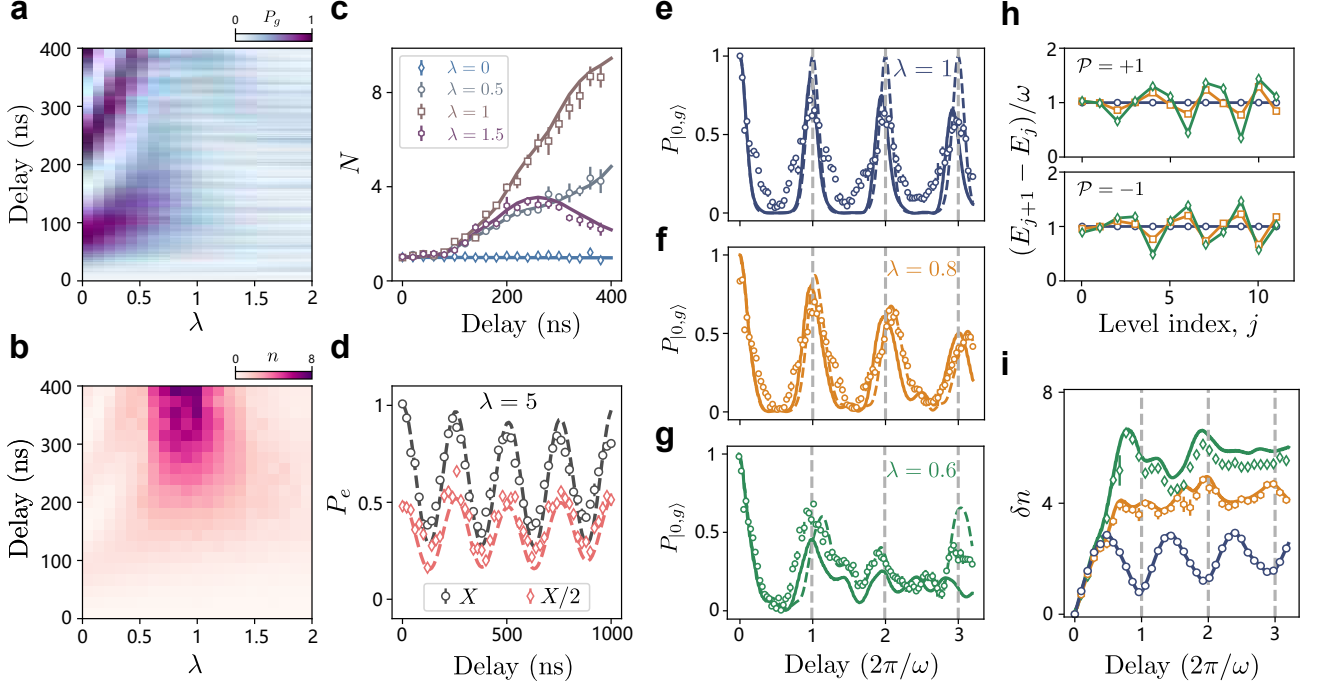


FIG. 2. **Tunable anisotropic interaction and collapse-revival dynamics in the AQRM.** **a,b,** Measured time-resolved qubit excitation probability P_g and cavity photon population n as functions of the anisotropy ratio λ . **c,** Extracted total photon number N for different values of λ . **d,** Evolution of the qubit populations for $\lambda = 5$ AQRM from the dual mapping, by applying X and $X/2$ gates to the system qubit respectively. The parameters are $\Omega/2\pi = \omega/2\pi = 1$ MHz and $g_0/2\pi = 2$ MHz. **e–g,** Collapse-revival dynamics for different values of λ in the degenerate case. The cases $\lambda = 1, 0.8$, and 0.6 are shown in blue, orange, and green, respectively. The parameters are $\omega/2\pi = 1.6$ MHz and $g_0/2\pi = 4.5$ MHz, corresponding to $\min\{g_1/\omega, g_2/\omega\} = 1.4, 1.25$, and 1.05 , respectively, which are all well within the DSC regime. **h,** Adjacent level spacings $(E_{j+1} - E_j)/\omega$ in the positive- and negative-parity sectors. **i,** Photon-number fluctuation δn extracted from the dynamics for different λ . Symbols represent experimental data, and dashed curves show numerical results from the effective AQRM. Solid curves include the full experimental protocol, decoherence, and the multilevel structure of the transmon treated as a qutrit.

Experimental setup

Our experiment is performed on a flip-chip superconducting quantum processor, comprising a central bus resonator R and 24 frequency-tunable transmon qubits [50]. We choose one transmon as the system qubit Q , coupled to R through a tunable coupler C for the in-situ control of the qubit–resonator interaction [51]. The intracavity photon state is measured using an ancillary transmon Q_a dispersively coupled to R , enabling photon-number-resolved detection [52]. The system qubit has an anharmonicity $\alpha/2\pi = -0.37$ GHz, which is sufficiently large to suppress leakage and support an effective two-level description.

The central capability of the experiment is to convert microwave control into the independent control of the two quadratures of the light–matter interaction. We engineer the AQRM by combining the longitudinal qubit frequency modulation $\tilde{\omega}_q$ with a multi-tone transverse drive $\mathcal{E}$. The laboratory-frame Hamiltonian is

$$\hat{H}_{\text{lab}} = \frac{\tilde{\omega}_q}{2} \hat{\sigma}_z + \omega \hat{a}^\dagger \hat{a} + (g_0 \hat{a}^\dagger \hat{\sigma}_- + \mathcal{E} \hat{\sigma}_- + \text{H.c.}), \quad (2)$$

where $\tilde{\omega}_q = \omega_q + A \cos(\mu t)$ and $\mathcal{E} = [V_0 + V \cos(\nu t)] e^{i\omega_0 t}$.

The mechanism can be understood from a spin-1/2 picture. The cavity field acts as a transverse magnetic field with two quadratures, $B^x = g_0(\hat{a} + \hat{a}^\dagger)/2$ and $B^y = ig_0(\hat{a}^\dagger - \hat{a})/2$ [27]. In the rotating frame with ω_0 , the carrier tone sets the bias as $\varepsilon = 2V_0$, while the transverse modulation produces an x -field, $B_0^x(t) = V \cos(\nu t)$. The rapid modulation of the total x field, $B^x + B_0^x(t)$, renormalizes the B^y component with the Bessel factor $J_0(2\beta)$, with $\beta = V/\nu$, because $\hat{\sigma}_y$ couples eigenstates of $\hat{\sigma}_x$. This mechanism is analogous to the dynamical control of hopping in the driven lattice systems [53–55]. Consequently, the two cavity quadratures acquire different effective strengths, and the stroboscopic dynamics are described by the AQRM Hamiltonian. The total coupling g_0 is partitioned into rotating and counterrotating components g_1 and g_2 , with $g_1 + g_2 = g_0$, and the anisotropy ratio becomes $\lambda = (1 - J_0(2\beta))/(1 + J_0(2\beta))$. This allows λ to be continuously tuned from 0 to 2.3 in our implementation. Because the effective USC/DSC coupling is generated with Floquet synthesis rather than through an intrinsically bare coupling, this approach avoids the equilibrium constraints associated with no-go arguments [56, 57].

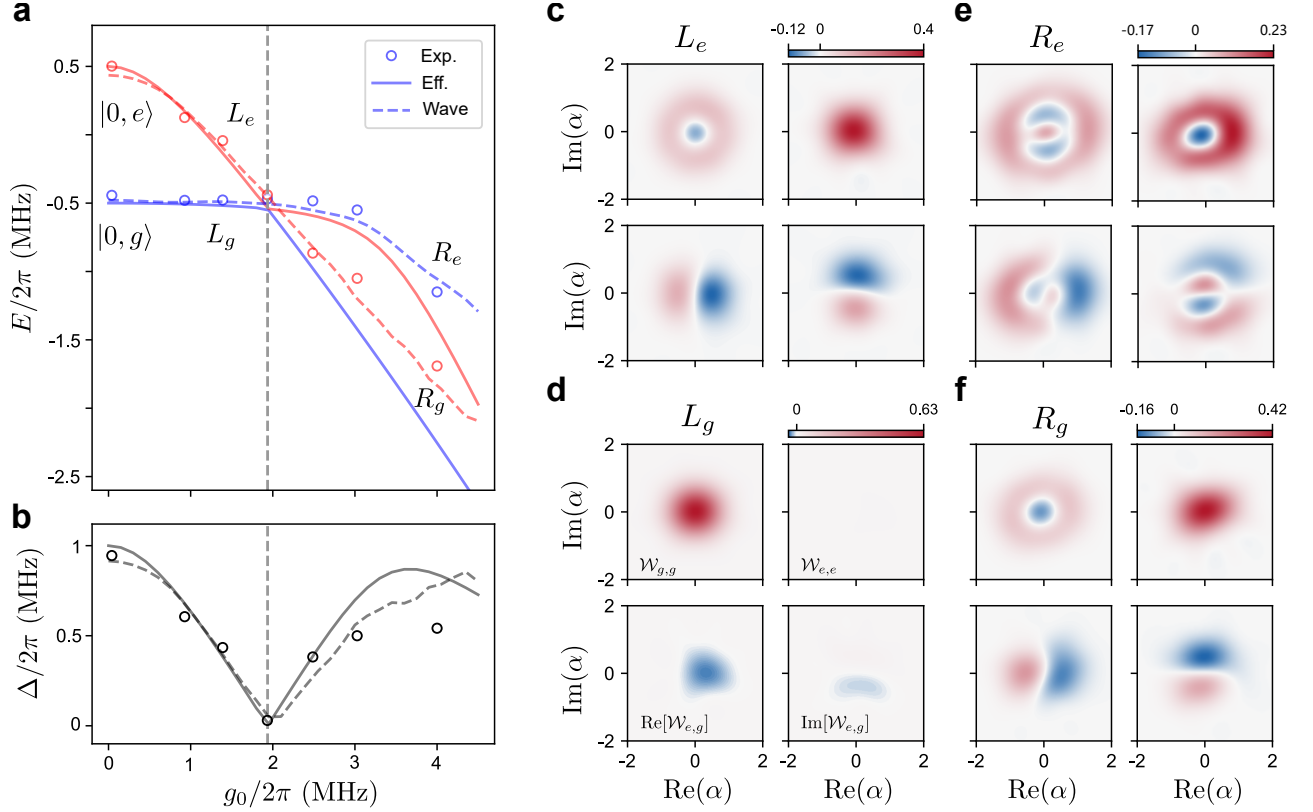


FIG. 3. **Ground-state level crossing and phase-space signatures of the AQRM.** **a**, Energy spectrum as a function of $g_0 = g_1 + g_2$, showing an exact crossing between the lowest even- and odd-parity branches. The solid curves correspond to numerical results calculated via the effective Hamiltonian, whereas the dashed curves are obtained using the same wave protocol adopted in the experiment. **b**, Energy gap between the ground state and the first excited state. The dashed line denotes the predicted crossing point g_{0c} . **c–f**, Joint Wigner function matrix at $g_0/2\pi = 1.5$ (4) MHz for $|0, g\rangle$, $|0, e\rangle$ as initial states, labelled as L_g (R_g) and L_e (R_e) individually. The parameters are $\Omega/2\pi = 1$ MHz, $\omega/2\pi = 2.5$ MHz, and $\lambda = 0.2$. The colorbars are set white when the value is zero, red when positive, and blue as negative.

We first verify that this control does more than rescale the interaction strength: it changes the symmetry of the dynamics. To benchmark the engineered interaction, we isolate $\hat{H}_{\text{int}} = g_1(\hat{H}_r + \lambda\hat{H}_{\text{cr}})$ and initialize the system in a single-photon state $|0, e\rangle$. Figures 2a and b show the measured qubit and resonator populations, as λ is swept at fixed $g_0/2\pi = 3$ MHz. We further show the extracted total photon number in Fig. 2c, which provides a compact probe of $U(1)$ -symmetry breaking and multiphoton generation. For $\lambda = 0$, N remains conserved, as expected from the JCM symmetry. Increasing λ activates counterrotating processes, mixes excitation manifolds, and drives the system from bounded JCM oscillations toward the photon growth of the isotropic point $\lambda = 1$, where $N \simeq g^2 t^2$ in the DSC limit [50, 58]. For $\lambda > 1$, the dynamics are suppressed, because $\hat{H}_{\text{cr}}$ does not act on $|0, e\rangle$ at first order. The extracted photon number captures this symmetry-controlled crossover and agrees well with theoretical prediction.

The rotating and counterrotating processes are further connected by introducing the identity $\hat{H}(\omega, \Omega, g_1, g_2, \varepsilon) =$

$\hat{\sigma}_x \hat{H}(\omega, -\Omega, g_2, g_1, \varepsilon) \hat{\sigma}_x$, which maps between $0 < \lambda < 1$ and $\lambda > 1$ sectors [50]. Experimentally, this duality allows a target AQRM with anisotropy λ to be realized by implementing its dual Hamiltonian at $1/\lambda$ with opposite qubit frequency, together with π pulses before and after the evolution. Using this scheme, we realize the dynamics at $\lambda = 5$ through the dual point $\lambda = 0.2$, for the initial states $|0, e\rangle$ and $(|0, g\rangle - i|0, e\rangle)/\sqrt{2}$, as shown in Fig. 2d. The agreement with the target dynamics demonstrates the access to the full AQRM parameter space, spanning the JCM, QRM, and anti-JCM limits [59].

Collapse-revival dynamics in the DSC regime

Collapse-revival dynamics provides a direct time-domain probe of the spectral regularity in the USC/DSC regime [19, 20] and has been widely explored in the QRM across diverse platforms [21–23, 60]. In Fig. 2e, we present the DSC dynamics of the QRM at $g/\omega = 1.4$ in the degenerate case $\Omega = 0$, initialized from the JCM ground state $|0, g\rangle$. The initial-state population almost decays to zero, remains suppressed over a period of $2\pi/\omega$, and then revives completely. This behaviour

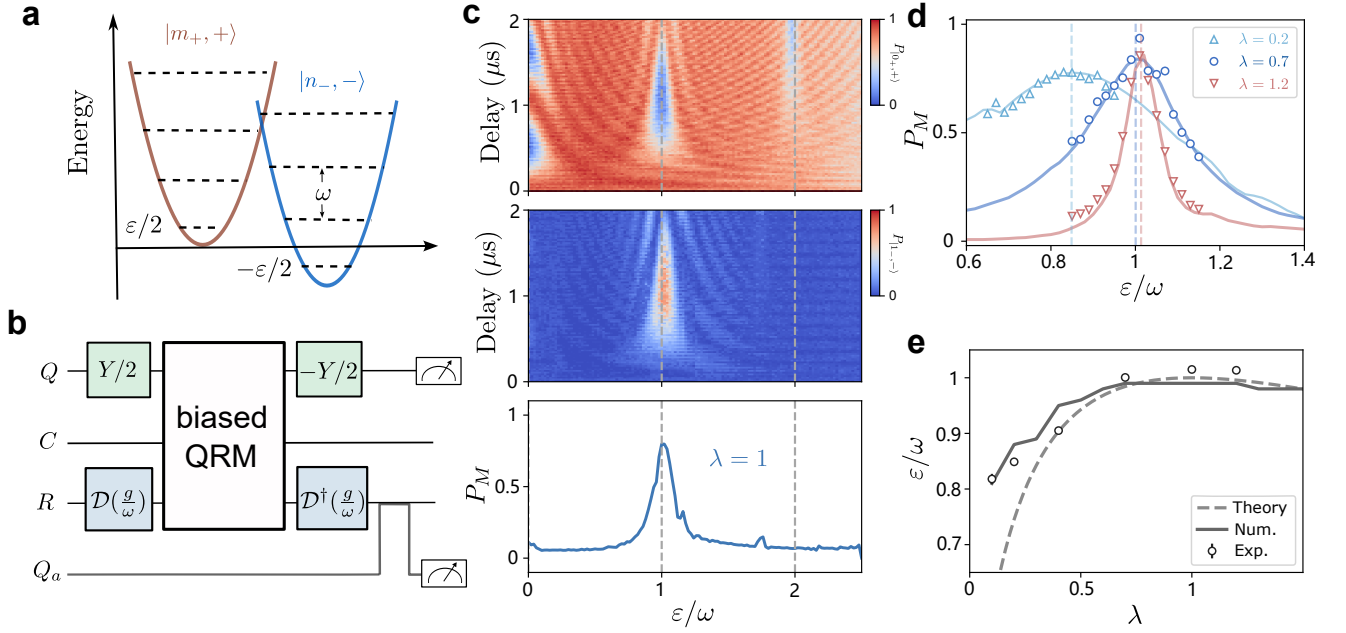


FIG. 4. **Hidden symmetry and tunnelling dynamics in biased models.** **a**, Schematic effective potentials of the biased QRM for $\varepsilon \neq 0$, where the $\mathbb{Z}_2$ symmetry is broken. **b**, Pulse sequence for observing tunnelling dynamics in $P_{|0+, +\rangle}$ in the biased QRM by mapping bare number states onto the displaced Fock states. **c**, The time evolution of QRM starting from the initial state $|0+, +\rangle$ against ε/ω . The upper, middle, and lower panels show $P_{|0+, +\rangle}$, $P_{|1-, -\rangle}$, and the extracted P_M , respectively. The dashed lines denote $\varepsilon/\omega = 0, 1, 2$. **d**, The transfer population P_M as a function of ε/ω for several values of λ . The dashed lines represent the position where P_M achieves a peak, where the solid curves denotes the numerical results. **e**, Extracted position of P_M peaks as a function of λ . The dashed curve denotes the theoretical hidden symmetry points, while the solid curve represent the numerical results using the same protocol and parameters as the experiments. Error bars represent one standard deviation.

can be understood from a displacement transformation $\mathcal{D}(g/\omega)$ [20], with $\mathcal{D}(\eta) = e^{\eta \hat{a}^\dagger - \eta^* \hat{a}}$. In the degenerate isotropic QRM, this transformation maps the system to an effective harmonic oscillator with equal level spacing ω , so the dynamical phases rephase exactly after one period. For $\lambda \neq 1$, however, the eigenenergies are roots of a transcendental function [30, 38] and are no longer equally spaced, interrupting the interference that produces perfect revivals.

In Figs. 2f and g, we present the DSC dynamics for the anisotropic cases for $\lambda = 0.8$ and 0.6 , respectively, using the same Ω , ω , and g_0 as in the isotropic case. The initial-state population still exhibits a collapse-revival pattern, but the revival is incomplete, which is consistent with the breakdown of the equally spaced spectral structure. In Fig. 2h, we show the corresponding level spacings $E_{j+1} - E_j$ in the two parity subspaces. We further characterize the dynamics using the photon-number fluctuation $\delta n = \sqrt{\langle \hat{n}^2 \rangle - \langle \hat{n} \rangle^2}$, as shown in Fig. 2i. For $\lambda = 1$, the state returns to $|0, g\rangle$ up to a global phase at integer multiples of $2\pi/\omega$ [20], yielding $\delta n = 0$ at those instants. For $\lambda = 0.6$ and 0.8 , the minima of δn are lifted, showing that anisotropy leaves a dynamical fingerprint even in the nominally degenerate limit. These results connect a programmable microscopic imbalance directly to a measurable reconstruction of the spectral interference.

Ground-state parity switching

Spectral reconstruction is also visible in the low-lying eigenstates. For $\varepsilon = 0$, the AQRM preserves the same $\mathbb{Z}_2$ parity symmetry as the QRM, splitting the spectrum into even- and odd-parity branches. The generic branches form the regular spectrum, while isolated crossings between opposite-parity branches occur only at the special parameter values, giving rise to the exceptional spectrum. A hallmark of anisotropy is an exact crossing between the two lowest levels of opposite parity. As g_0 increases at a fixed λ , this crossing exchanges the ground and first-excited states, producing a ground-state parity switch absent in the isotropic QRM. For the first crossing with $0 \leq \lambda < 1$, the crossing point is given by the Judd-type exceptional solution $g_{0c} = (1 + \lambda)g_{1c}$ [30], where $g_{1c} = \sqrt{\Omega\omega/(1 - \lambda^2)}$. This provides a sharp spectral landmark of the AQRM.

To locate this crossing experimentally, we track the two lowest energies using adiabatic state preparation followed by joint qubit-resonator tomography [52]. Starting from the decoupled point, the qubit is initialized in $|g\rangle$ or $|e\rangle$ and the coupler is adiabatically ramped to the target interaction strength, mapping the product states onto low-lying eigenstates of the coupled Hamiltonian. Tomography yields the joint density matrix ρ and the energy expectation $\text{Tr}[\rho \hat{H}]$, allowing the two branches and

their interchange to be resolved. In Figs. 3a and b, we measure the ground and first-excited energies as a function of g_0 and identify an exact crossing at the predicted g_{0c} .

The joint qubit-resonator state can be fully characterized by a Wigner function matrix [27, 52, 61–64], with elements $\mathcal{W}_{kk'}(\alpha)$ constructed from the cavity blocks $\rho_{kk'} = \langle k|\rho|k'\rangle$ for $k, k' \in \{g, e\}$. Since the Wigner value at the origin gives the photon-number parity $\mathcal{W}(0) = 2\langle(-1)^{\hat{n}}\rangle/\pi$ [5, 61, 62], the total parity is obtained from the qubit-resolved components as

$$\mathcal{P} = \frac{\pi}{2} [\mathcal{W}_{gg}(0) - \mathcal{W}_{ee}(0)]. \quad (3)$$

The measured Wigner function matrices in Fig. 3c-f show that the ground state changes from even parity ($\mathcal{P} = 1.02$ for $g_0 < g_{0c}$) to odd parity ($\mathcal{P} = -0.91$ for $g_0 > g_{0c}$). This directly reveals an anisotropy-induced parity switch from the competition between rotating and counterrotating channels.

Hidden symmetry and tunnelling dynamics

A final test of the programmable symmetry is provided by the biased model. Although a nonzero bias breaks the $\mathbb{Z}_2$ parity symmetry and generally lifts parity-protected degeneracies, exact level crossings can reemerge at discrete commensurate biases owing to hidden symmetries. At these points, wave-packet dynamics are confined to a reduced subspace spanned by displaced polaron states, enabling coherent tunnelling between quasi-degenerate effective potential wells [42], as illustrated in Fig. 4a. For the QRM, the hidden symmetry emerges when the bias satisfies the commensurate condition

$$\varepsilon_h/\omega = k, \quad (4)$$

where k is a positive integer. Under this resonance condition, the tunnelling is mediated by states that are direct tensor products of displaced Fock states and transversely polarized qubit states, $|n_{\pm}, \pm\rangle = \mathcal{D}(\pm g/\omega)|n\rangle \otimes |\pm\rangle$, where $|\pm\rangle = (|e\rangle \pm |g\rangle)/\sqrt{2}$. Because this tunnelling is a high-order, symmetry-selected resonance with a narrow linewidth, it is sensitive to parameter fluctuations and pulse imperfections, making its observation a stringent test of the simulator control accuracy. Experimentally, we align the evolution, initial state, and measurement basis with the hidden-symmetry subspace using joint displacement operations and single-qubit gates before and after the biased QRM's evolution, as shown in Fig. 4b. These operations map the displaced Fock-state basis onto the bare photon-number basis for measurement. Figure. 4c shows the dynamics of $P_{|0_{+}, +\rangle}$ and $P_{|1_{-}, -\rangle}$ as a function of ε/ω , initialized from $|0_{+}, +\rangle$. In the lower panel, we plot the extracted maximum transfer probability, $P_M = \max_t P_{|1_{-}, -\rangle}(t)$, as a function of ε/ω . The tunnelling points originated from hidden symmetries are clearly identified near $\varepsilon/\omega = 1.0$, corresponding to the commensurate biases in Eq. (4).

For $\lambda \neq 1$, the hidden-symmetry points become λ dependent and are given by [42]

$$\varepsilon_h/\omega = k \frac{2\sqrt{\lambda}}{1+\lambda}, \quad (5)$$

which reduces to Eq. (4) at $\lambda = 1$. Although an analogous tunnelling-dynamics picture can be intuitively constructed, obtaining an analytic form of the corresponding tunnelling initial state is difficult in the anisotropic case. We therefore prepare a λ -dressed initial state, $|0_{+}, +\rangle_{\lambda} = \mathcal{D}(-\sqrt{\lambda}g_1/\omega)|0\rangle \otimes (\sqrt{\lambda}|g\rangle + |e\rangle)/\sqrt{1+\lambda}$ and measure the population at $|1_{-}, -\rangle_{\lambda} = \mathcal{D}(\sqrt{\lambda}g_1/\omega)|1\rangle \otimes (\sqrt{\lambda}|g\rangle - \sqrt{\lambda}|e\rangle)/\sqrt{1+\lambda}$, which can be effectively derived in the weak-coupling regime [50]. Figure 4d shows P_M as a function of ε/ω , revealing a λ -dependent shift of the peak positions. In Fig. 4e, we extract the value of ε/ω at which P_M reaches its maximum and plot it as a function of λ . The extracted trend follows the predicted anisotropy-dependent displacement of the hidden-symmetry point in Eq. (5), providing a dynamical signature of generalized hidden symmetry in the AQRM.

Conclusion and outlook

In summary, we have realized a programmable extension of the quantum Rabi model in which anisotropy and bias become experimentally tunable axes of symmetry control. This capability goes beyond reaching a large effective coupling: it allows the balance between energy-conserving and nonconserving processes to be engineered, thereby controlling spectral structure, protected crossings and dynamical selection rules. The observed incomplete revivals, ground-state parity switching and hidden-symmetry tunnelling provide complementary views of the same underlying principle: the nonperturbative light-matter dynamics can be shaped by programming the symmetry of the Hamiltonian. Looking forward, the approach can be extended to time-dependent anisotropic protocols, multi-qubit anisotropic Dicke-type models [28, 65], and more complex Floquet-engineered light-matter Hamiltonians [66].

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Code availability

The codes used for the numerical simulations are available from the corresponding authors upon request.

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Competing interests

The authors declare no competing interests.

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