Two-level system hyperpolarization using a quantum Szilard engine

Martin Spiecker 1,2, Patrick Paluch 1,2, Nicolas Gosling 2, Niv Drucker 3, Shlomi Matityahu 4, Daria Gusenko 1,2, Simon Günzler 1,2, Dennis Rieger 1,2, Ivan Takmakov 1,2, Francesco Valenti 2, Patrick Winkel 1,2, Richard Gebauer 5, Oliver Sander 5, Gianluigi Catelani 6, Alexander Shnirman 2,4, Alexey V. Ustinov 1,2, Wolfgang Wernsdorfer 1,2, Yonatan Cohen 3 & Ioan M. Pop 1,2

The innate complexity of solid-state physics exposes superconducting quantum circuits to interactions with uncontrolled degrees of freedom degrading their coherence. By implementing a quantum Szilard engine with an active feedback control loop, we show that a superconducting fluxonium qubit is coupled to a two-level system (TLS) environment of unknown origin, with a relatively long intrinsic energy relaxation time exceeding 50 ms. The TLSs can be cooled down, resulting in a four times lower qubit population, or they can be heated to manifest themselves as a negative-temperature environment corresponding to a qubit population of ~80%. We show that the TLSs and qubit are the dominant loss mechanism for each other and that qubit relaxation is independent of the TLS populations. Understanding and mitigating TLS environments is, therefore, not only crucial to improve the qubit lifetimes but also to avoid non-Markovian qubit dynamics.

Although tremendous progress has been made to improve the coherence of superconducting qubits, they still have to cope with various loss and decoherence mechanisms, certainly to the chagrin of quantum computing scientists but also to the joy of mesoscopic physicists. The relentless interactions between superconducting hardware and its environment motivate the development of quantum error correction using stabilizer codes on one hand1–4 and deepen our understanding of mesoscopic processes on the other hand5–15. In the past, numerous strategies have been conceived to study and mitigate decoherence from various sources, from defects in dielectrics to non-thermal excitations16. A major source can be attributed to the wide class of two-level system (TLSs) in the qubit environment. Weakly coupled TLSs may be investigated by saturation pulses17,18, whereas strongly coupled TLSs may even be coherently operated via the superconducting qubit19,20. Moreover, it has been shown that a sequence of repeated π pulses can change the environment of the superconducting qubit, which was interpreted as the diffusion of superconducting quasiparticles away from the qubit junctions9.

Here we implement a quantum Szilard engine21–24 that manipulates the environment of a superconducting qubit. Our Szilard engine executes a hyperpolarization protocol similar to experiments using spin qubits25 or defect centres26, and is readily applicable in state-of-the-art quantum processors27–29. The hyperpolarized environment reveals that the qubit is weakly coupled to a TLS environment of unknown origin, which relaxes over tens of milliseconds. Conversely, this previously hidden environment can now be identified as the dominant loss mechanism of our qubit, and we dread that similarly acting environments are ubiquitous in superconducting hardware. The quantum Szilard engine consists of a granular aluminium fluxonium qubit30 that can be actively prepared in one of its eigenstates |g⟩ or |e⟩. The fluxonium and its complex environment are depicted in Fig. 1a. The Szilard engine implements a dynamical polarization protocol on the TLSs (Fig. 1b,c). In contrast to

1PHI, Karlsruhe Institute of Technology, Karlsruhe, Germany. 2IQMT, Karlsruhe Institute of Technology, Eggenstein-Leopoldshafen, Germany. 3Quantum Machines, Tel Aviv-Yafo, Israel. 4TKM, Karlsruhe Institute of Technology, Karlsruhe, Germany. 5IPE, Karlsruhe Institute of Technology, Eggenstein-Leopoldshafen, Germany. 6JARA Institute for Quantum Information (PGI-11), Forschungszentrum Jülich, Jülich, Germany.

*e-mail: martin.spiecker@kit.edu; ioan.pop@kit.edu
The cross-relaxation rates $\Gamma_\uparrow = 3.4\, \text{‰}$, respectively, capturing the quasiparticles, absorbed molecules on the surface\(^{52}\) and dielectric TLSs\(^{53}\). The relaxation of the reservoir cannot be directly observed and has to be inferred from the qubit dynamics. Although the common approach is to measure the free decay of the qubit (Supplementary Section A), here we exploit the fact that the qubit readout is more than 96% quantum non-demolishing (Supplementary Section B and ref. 31) and we perform repeated single-shot readouts, resulting in stroboscopic quantum-jump traces (Fig. 2d). The constant relaxation rate $\Gamma_q \approx 10^3$, during stabilization in $|g\rangle$ or $|e\rangle$, respectively. The time-evolving transition rates (Fig. 4a) are obtained from the rate equations read as follows:

$$\Gamma_k = \Gamma_{\uparrow} + \Gamma_{\downarrow}$$

and the equilibrium population of the qubit $p_{eq} = \Gamma_{\downarrow}/\Gamma_{\uparrow}$. Note that the noise magnitude varies with the qubit population (Fig. 4a), because the rates $\Gamma_{\downarrow}$ are based on conditional probabilities. Remarkably, after a heating sequence with $N = 10^4$, $\Gamma_{\uparrow}$ of the qubit is comparably constant (Fig. 4b); in contrast, $\Gamma_{\downarrow}$ follows a non-exponential relaxation for timescales up to 50 ms. At the end of the polarization sequence, for the TLSs, we can ascribe a hyperpolarization $p_{eq}^{TLS} = 97\%$, which—when taking into account the intrinsic loss of the qubit—gives the measured $p_{eq} = 78\%$ (Fig. 4c). Conversely, after a cooling sequence with $N = 10^4$, we extract $p_{eq} = 3.0\%$, as can be ascertained in Fig. 3b using the qubit population $p_e = 2.0\% = p_{eq}$ after $\Gamma_{\downarrow}$. Hence, the Szilard engine cooled the environment to an effective temperature of 16 mK, which is well below the temperature of the dilution refrigerator (~25 mK) and the effective temperature $T_{eq} = 28.3$ mK corresponding to the idle qubit population $p_{eq} = 12.0\%$ (Fig. 2d). The TLS hyperpolarization is even lower, $p_{eq}^{TLS} = 3.4\% \pm 9.9$ mK, limited by the qubit preparation infidelity. The values are extrapolated from the theoretical model, which will be explained in the next paragraph. For both heating and cooling, the hyperpolarization values are among the highest reported in the literature\(^{20,31}\).

The constant relaxation rate $\Gamma_q$ as well as the observed population inversion indicate an environment consisting of TLSs. We, therefore, model the system assuming the qubit to be coupled to a countable number of TLSs with populations $p_k^{TLS}$. The cross-relaxation rates $I_{qt}^k$ between the qubit and TLSs are given by\(^{34}\)

$$I_{qt}^k = \frac{2 \delta_k}{\Gamma_q^k}$$

where $\delta_k$ is the detuning between the qubit and $k$th TLS, $g$ is their transverse coupling strength and $\Gamma_q^k$ is the sum of their decoherence rates. Since the TLSs can, in turn, excite the qubit, we conclude that the qubit and TLSs are close in frequency so that they approximately relax to the same thermal population $p_{th}$ (note that the qubit is well thermalized, as discussed earlier). Finally, we introduce intrinsic relaxation rates for the qubit and TLSs, namely, $\Gamma_q^1$ and $\Gamma_{\uparrow}^k$, respectively, capturing the remaining environment (Fig. 1b). The dynamics is governed by the so-called Solomon equations\(^{40}\), extensively used in the field of nuclear hyperpolarization\(^{35}\). The rate equations read as follows:
Parametric Amplifier; RO, readout. \( \Gamma \) is consistent with the fact that we do not show in (Supplementary Section D provides a schematic of the detailed setup). We start the experiment with the TLS polarization sequence (Fig. 1c) by stabilizing the qubit to either \( |g⟩ \) or \( |e⟩ \) using V active feedback preparations. This is followed by a qubit initialization in \( |g⟩ \) or \( |e⟩ \); immediately thereafter, we begin to stroboscopically monitor the qubit state. For polarization and qubit monitoring, the repetition time is \( t_{\text{rep}} = 2 \mu s \), much shorter than the qubit relaxation time \( T_{1} = 20 \mu s \). Before each of the 2,500 repetitions, we wait for 50 ms to allow the environment to relax. The protocol is orchestrated by the field-programmable gate array (FPGA) controller from Quantum Machines, with an internal real-time feedback latency of \( \sim 200 \) ns (Supplementary Section E provides a schematic of the detailed setup). \( a \). Scatter plot of the complex reflection coefficient \( S_{00} \) of the readout signal for the qubit in equilibrium (left) and after \( |e⟩ \) - state preparation (right). The readout integration time is \( 128 \) ns, resulting in a separation of \( 5.6 \) fs (green circles indicate \( 2 \).) \( b \). Typical quantum-jump trace during qubit monitoring (as shown in \( a \)). The solid line indicates the assigned qubit state. JPA, Josephson Parametric Amplifier; RO, readout.

\[
\dot{p}_{q} = -\Gamma_{1}(p_{q} - p_{\text{th}}) - \sum_{k} \Gamma_{k}^{\text{eq}}(p_{k} - p_{q}^{0}),
\]

\[
\dot{p}_{q}^{k} = -\Gamma_{k}^{\text{eq}}(p_{k} - p_{\text{th}}) - \Gamma_{k}^{\text{ab}}(p_{q}^{0} - p_{q}).
\]

where we identify the constant qubit relaxation rate as \( \Gamma_{1} = \Gamma_{g} + \sum_{k} \Gamma_{k}^{\text{eq}} \) and the time-dependent \( p_{\text{eq}} = (\Gamma_{1} p_{\text{th}} + \sum_{k} \Gamma_{k}^{\text{ab}} p_{k}^{0}) / \Gamma_{1} \). As a consequence of equation (3), during the polarization time \( N \times t_{\text{rep}} \), when we enforce \( p_{k} = 0 \) or 1, there is an exponential population transfer between the qubit and each TLS: at the end of the sequence, we expect to find the TLSs polarized (Fig. 2c).

So far, the model in equations (2) and (3) requires two rates for each TLS. To extract meaningful information from the measurements by virtue of equation (1), we need to make simplifying assumptions and reduce the number of fitting parameters. Since we observe the TLS polarization in different qubits and at different qubit frequencies (Supplementary Section I), we expect the TLSs to be randomly distributed in frequency. We simplify this distribution by modeling them to be equally spaced in frequency with \( \Delta_{0} = k \delta + \Delta_{\alpha} \), where \( \Delta_{\alpha} \in (0, \Delta/2) \) defines a shift in the TLS ladder with respect to the qubit frequency. This is justified by the fact that we are mainly interested in capturing the slow, non-exponential relaxation at millisecond timescales. With the same argument for all TLSs, we assume the same \( g_{k} \) and \( \Gamma_{k} \). The price we pay for using these simplifications is that the model less accurately captures the initial features of the decay curves, at \( t < 300 \) ms. Indeed, these features are a fingerprint of the exact configuration of the TLSs; as expected, they fluctuate in time \( ^{[16,17]} \) (Supplementary Section A).

The simplified model allows to rewrite equation (1) in the compact form \( I_{q}^{k} = g_{k}^{2} / | \beta_{k}^{\text{eq}} - (k + Bc_{k}) |^{2} \), showing that \( g_{k}, \Delta_{\alpha}, \) and \( \Gamma_{k} \) do not independently appear in the model. Instead, \( g = \sqrt{\text{diag}2/2} \Delta = \Gamma_{b}/2 \) and \( \Delta_{\alpha} = c_{f}^{2} \) can be determined for a given decoherence rate from a successful fit of the model. The fit procedure is further restricted by inserting the measured qubit relaxation rate as \( \Gamma_{q} = 1/21.5 \) (Fig. 4b), leaving us with only two essential fit parameters, namely, \( \Gamma_{q} \) and \( b \) (Supplementary Section J). The robustness of the model is illustrated by the fact that a fit of only the first millisecond to one of the stronger polarized relaxation curves (for example, polarization to \( |e⟩ \)) is sufficient to describe the highly non-exponential relaxation of all the measurements on the entire relaxation range up to 50 ms (Figs. 3 and 4, continuous lines). Details of the fitting procedure are presented in Supplementary Section J.

Using the lower bound of \( \Gamma_{2} \geq \Gamma_{q}^{2} = 0.5 \) MHz, where \( \Gamma_{q}^{2} \) is the decoherence rate of the qubit (Supplementary Section K), we extract \( g \geq 2n \times 12 \) kHz and \( \Delta \geq 2n \times 167 \) kHz. The comparably small coupling strength \( g \ll \Gamma_{q}^{2} \) is consistent with the fact that we do not observe avoided level crossings in the qubit spectrum. In particular, this argument remains valid even for higher decoherence because \( g \) and \( \Delta \) scale with \( \sqrt{T_{1}} \) and \( \Gamma_{2} \), respectively. Using an upper bound for the decoherence as \( \Gamma_{q} = 1/10n \) MHz, comparable with values reported in another work \( ^{[29]} \), gives \( g < 2 \pi \times 170 \) kHz and \( \Delta < 2 \times 35 \) MHz.

Furthermore, we can calculate the two contributions of the qubit relaxation: one rate is due to interactions with the TLSs, that is, \( r_{\text{TLS}} = \sum_{k} r_{q}^{k} = 35.9 \) kHz, and the other is the remaining intrinsic relaxation \( T_{\text{rel}} = 10.7 \) kHz. We, therefore, identify the TLS bath as the dominant loss mechanism. Remarkably, the fit also indicates that the intrinsic relaxation time exceeds \( 1/\Gamma_{2} > 50 \) ms, which is orders of magnitude longer than previously measured relaxation rates of dielectric TLSs \( ^{[30,40]} \). This fact leads us to believe that we are reporting a new type of TLS environment, possibly related to spins \( ^{[41]} \) or trapped quasiparticle TLSs \( ^{[42]} \). Finally, we would like to mention that \( r_{q}^{k} \geq \Gamma_{q}^{2} / |k| \leq 15 \), which means that the qubit is the main decay channel for at least the first few tens of the most resonant TLSs.

Following Szilard’s seminal paper \( ^{[22]} \), the homonymous engine uses measured information as fuel (Supplementary Section L). In the first iteration of a cooling sequence starting from thermal equilibrium at \( T = 28.3 \) mK, the engine extracts, on average, the internal energy as fuel (Supplementary Section A), \( U_{\text{out}} = 0.24 k_{B} T \) from the qubit, corresponding to an entropy reduction of \( 0.37 k_{B} \), which should be compared with the entropy produced by the measurement apparatus \( k_{B} \ln 2 = 0.69 k_{B} \). From the rate equation, we can calculate the optimal working regime for our Szilard engine. Using the fitted parameters, we infer that the maximum
heat reduction of $\Delta Q = 0.11k_B T$ in the reservoir occurs 68 $\mu$s after qubit initialization. Thus, at most half of the extracted heat from the qubit can be used to cool the reservoir. With a similar timescale of $t_{\text{rep}} = 100 \mu$s (Supplementary Section M), we show that the reservoir can also be heated by a sequence of $\pi$ pulses. However, this procedure, introduced in another work\textsuperscript{9}, cannot result in population inversion due to interactions with the environment. In particular, the $|f\rangle$ state is not populated, as illustrated by the absence of a third cloud in the $S_{11}$ distribution (Supplementary Section H).

In summary, using a superconducting qubit and active feedback, we demonstrated a quantum Szilard engine that can polarize a TLS environment of unknown origin. As a result, the qubit population exhibits remarkably long and non-exponential dynamics due to the intrinsically long decay time of the TLSs, exceeding 50 ms. This showcases the challenges and pitfalls of extracting $T_1$ from the relaxation data of the qubit population. In our device, we extract $T_1$ from quantum jumps and show that it is unaffected by the continuous operation of the qubit, ruling out enhanced quasiparticle diffusion\textsuperscript{9}. Although $T_1$ is independent of the environment population, the transition rates $T_{\text{eff}}$ are not. Our results are particularly relevant in the context of quantum processors, where the heating and cooling of the environment is a byproduct of continuous operation. The Szilard engine could be used to study out-of-equilibrium processes or to preferentially reduce one of the qubit transition rates. For example, reducing $T_{\text{eff}}$ would be beneficial for bosonic codes\textsuperscript{43,44}.

In our system, quantum coherence between the qubit and TLSs can be neglected, allowing a simple description using the Solomon equations. As quantum hardware continues to improve, coherent interactions and non-Markovian qubit dynamics will start to play a role, raising the bar for quantum error correction strategies. The quantum Szilard engine presented here offers the first glimpse of the challenges facing future hardware, in which coherence improvements also translate into increasingly complex interactions with the environment.
Fig. 4 | Measured population inversion and constant relaxation rate—signatures of the TLS environment. a. Measured (light colour) and calculated (dark colour) qubit transitions rates $\Gamma_{\uparrow,\downarrow}$ following initialization in |g⟩ or |e⟩ and for increasing polarization to |e⟩ from $N = 10$ to $10^3$ and $10^4$ (left, middle and right panels, respectively). The measured rates are extracted from the same quantum-jump traces used to extract the qubit relaxation (Fig. 3a), with the logarithmic time axis starting at 500 μs. For long polarization times, the rates are reversed in the beginning, meaning that the qubit sees a negative-temperature environment. Note that in all the cases, the |g⟩-state initialization visibly cools the environment, suggesting a heat capacity of only a few energy quanta. To reduce the statistical noise, a five-point moving-average filter was applied corresponding to a 10 μs window. Furthermore, the first 10 μs of the orange and green curves are omitted due to low statistics; it is unlikely to have two consecutive intervals of ≤5 μs between jumps. Similarly, these rates are overestimated in the beginning by preferably detecting short $\Gamma$ fluctuations of the qubit. b. Relaxation time $\Gamma_1$ obtained from the $\Gamma_{\uparrow,\downarrow}$ rates in a for $N = 10^4$. The $\Gamma_1$ rate is constant compared with the changes in $\Gamma_\uparrow$ and $\Gamma_\downarrow$ (shown in a), that is, $\Gamma_1$ is independent of the environmental populations, which indicates a TLS environment. c. Equilibrium population of the qubit $\rho_{eq}$ extracted from the $\Gamma_{\uparrow,\downarrow}$ rates. The dashed lines show the corresponding qubit population $\rho_{eq}$ relaxation taken from Fig. 3a. We extrapolate an effective population of the environment $\rho_{eq}$ = 78% at $t = 0$. In all panels, the theoretical curves use the same parameters as that in Fig. 3.

References
1. Ofek, N. et al. Extending the lifetime of a quantum bit with error correction in superconducting circuits. Nature 536, 441–445 (2016).
2. Vuillot, C., Asasi, H., Wang, Y., Pryadko, L. P. & Terhal, B. M. Quantum error correction with the toric Gottesman-Kitaev-Preskill code. Phys. Rev. A 99, 032344 (2019).
3. Google Quantum AI. Exponential suppression of bit or phase errors with cyclic error correction. Nature 595, 383–387 (2021).
4. Cai, W., Ma, Y., Wang, W., Zou, C.-L. & Sun, L. Bosonic quantum error correction codes in superconducting quantum circuits. Fundam. Res. 1, 50–67 (2021).
5. Grabovskij, G. J., Peichl, T., Lisenfeld, J., Weiss, G. & Ustinov, A. V. Strain tuning of individual atomic tunneling systems detected by a superconducting qubit. Science 338, 232–234 (2012).
6. Barends, R. et al. Coherent Josephson qubit suitable for scalable quantum integrated circuits. Phys. Rev. Lett. 111, 080502 (2013).
7. Risté, D. et al. Millisecond charge-parity fluctuations and induced decoherence in a superconducting transmon qubit. Nat. Commun. 4, 1913 (2013).
8. Pop, I. M. et al. Coherent suppression of electromagnetic dissipation due to superconducting quasiparticles. Nature 508, 369–372 (2014).
9. Gustavsson, S. et al. Suppressing relaxation in superconducting qubits by quasiparticle pumping. Science 354, 1573–1577 (2016).
10. Grünhaupt, L. et al. Loss mechanisms and quasiparticle dynamics in superconducting microwave resonators made of thin-film granular aluminum. Phys. Rev. Lett. 121, 117001 (2018).
11. Serniak, K. et al. Hot nonequilibrium quasiparticles in transmon qubits. Phys. Rev. Lett. 121, 157701 (2018).
12. Chu, Y. et al. Creation and control of multi-phonon Fock states in a bulk acoustic-wave resonator. Nature 563, 666–670 (2018).
13. De Graaf, S. E. et al. Two-level systems in superconducting quantum devices due to trapped quasiparticles. Sci. Adv. 6, eabc5055 (2020).
14. Wilen, C. D. et al. Correlated charge noise and relaxation errors in superconducting qubits. Nature 594, 369–373 (2021).
15. Glazman, L. & Catelani, G. Bogoliubov quasiparticles in superconducting qubits. SciPost Phys. Lect. Notes 31 (2021).
16. Siddiqi, I. Engineering high-coherence superconducting qubits. Nat. Rev. Mater. 6, 875–891 (2021).
17. Kirsh, N., Svetitsky, E., Burin, A. L., Schechter, M. & Katz, N. Revealing the nonlinear response of a tunneling two-level system ensemble using coupled modes. Phys. Rev. Mater. 1, 012601 (2017).
18. Andersson, G. et al. Acoustic spectral hole-burning in a two-level system ensemble. npj Quantum Inf. 7, 15 (2021).
19. Wang, Z. L. et al. Quantum state characterization of a fast tunable superconducting resonator. Appl. Phys. Lett. 102, 163503 (2013).
20. Lisenfeld, J. et al. Electric field spectroscopy of material defects in transmon qubits. npj Quantum Inf. 5, 105 (2019).
21. Szilard, L. Über die Entropieverminderung in einem thermodynamischen System bei Eingriffen intelligenter Wesen. Z. Phys. 53, 840–856 (1929).
22. Toyabe, S., Sagawa, T., Ueda, M., Muneyuki, E. & Sano, M. Experimental demonstration of information-to-energy conversion and validation of the generalized Jarzynski equality. Nat. Phys. 6, 988–992 (2010).
23. Koski, J. V., Maisi, V. F., Pekola, J. P. & Averin, D. V. Experimental realization of a Szilard engine with a single electron. Proc. Natl Acad. Sci. USA 111, 13786–13789 (2014).
24. Peterson, J. P. S., Sarthour, R. S. & Laflamme, R. Implementation of a quantum engine fuelled by information. Preprint at https://arxiv.org/abs/2006.10136 (2020).
25. Bluhm, H., Foletti, S., Mahalu, D., Umansky, V. & Yacoby, A. Enhancing the coherence of a spin qubit by operating it as a feedback loop that controls its nuclear spin bath. Phys. Rev. Lett. 105, 216803 (2010).
26. Broadway, D. A. et al. Quantum probe hyperpolarisation of molecular nuclear spins. Nat. Commun. 9, 1246 (2018).
27. Córcoles, A. D. et al. Exploiting dynamic quantum circuits in a quantum algorithm with superconducting qubits. Phys. Rev. Lett. 127, 100501 (2021).
28. Gold, A. et al. Entanglement across separate silicon dies in a modular superconducting qubit device. npj Quantum Inf. 7, 142 (2021).
29. Satzinger, K. J. et al. Realizing topologically ordered states on a quantum processor. Science 374, 1237–1241 (2021).
30. Grünhaupt, L. et al. Granular aluminium as a superconducting material for high-impedance quantum circuits. Nat. Mater. 18, 816–819 (2019).
31. Gusenkova, D. et al. Quantum nondemolition dispersive readout of a superconducting artificial atom using large photon numbers. Phys. Rev. Applied 15, 064030 (2021).
32. Yang, A. et al. Simultaneous subsecond hyperpolarization of the nuclear and electron spins of phosphorus in silicon by optical pumping of exciton transitions. Phys. Rev. Lett. 102, 257401 (2009).
33. Wang, H.-J. et al. Sensitive magnetic control of ensemble nuclear spin hyperpolarization in diamond. Nat. Commun. 4, 1940 (2013).
34. Solomon, I. Relaxation processes in a system of two spins. Phys. Rev. 99, 559–565 (1955).
35. Vögelü, B. The nuclear Overhauser effect from a quantitative perspective. Prog. Nucl. Magn. Reson. Spectrosc. 78, 1–46 (2014).
36. Klimov, P. V. et al. Fluctuations of energy-relaxation times in superconducting qubits. Phys. Rev. Lett. 121, 090502 (2018).
37. Thorbeck, T., Eddins, A., Lauer, I., McClure, D. T. & Carroll, M. TLS dynamics in a superconducting qubit due to background ionizing radiation. Preprint at https://arxiv.org/abs/2210.04780 (2022).
38. Neeley, M. et al. Process tomography of quantum memory in a Josephson-phase qubit coupled to a two-level state. Nat. Phys. 4, 523–526 (2008).
39. Lisenfeld, J. et al. Measuring the temperature dependence of individual two-level systems by direct coherent control. Phys. Rev. Lett. 105, 230504 (2010).
40. Lisenfeld, J. et al. Decoherence spectroscopy with individual two-level tunneling defects. Sci. Rep. 6, 23786 (2016).
41. Lee, K. H., Holmberg, O. E. & Crawford, J. H. Optical and ESR studies of hole centers in γ-irradiated Al2O3. Phys. Stat. Sol. A 39, 669–674 (1977).
42. Yang, F. et al. Microscopic charging and in-gap states in superconducting granular aluminum. Phys. Rev. B 102, 104502 (2020).
43. Reinhold, P. et al. Error-corrected gates on an encoded qubit. Nat. Phys. 16, 822–826 (2020).
44. Grimm, A. et al. Stabilization and operation of a Kerr-cat qubit. Nature 584, 205–209 (2020).
45. Samkharadze, N. et al. High-kinetic-inductance superconducting nanowire resonators for circuit QED in a magnetic field. Phys. Rev. Applied 5, 044004 (2020).
46. Borisov, K. et al. Superconducting granular aluminum resonators resilient to magnetic fields up to 1 Tesla. Appl. Phys. Lett. 117, 120502 (2020).
47. Quintana, C. M. et al. Observation of classical-quantum crossover of 1/f flux noise and its paramagnetic temperature dependence. Phys. Rev. Lett. 118, 057702 (2017).
48. Reed, M. D. et al. Fast reset and suppressing spontaneous emission of a superconducting qubit. Appl. Phys. Lett. 96, 203110 (2010).
49. von Lüpke, U. et al. Parity measurement in the strong dispersive regime of circuit quantum acoustodynamics. Nat. Phys. 18, 794–799 (2022).
50. Vool, U. et al. Non-Poissonian quantum jumps of a fluxonium qubit due to quasiparticle excitations. Phys. Rev. Lett. 113, 247001 (2014).
51. Nsanzinez, I. & Plourde, B. L. T. Trapping a single vortex and reducing quasiparticles in a superconducting resonator. Phys. Rev. Lett. 113, 117002 (2014).
52. Kumar, P. et al. Origin and reduction of 1/f magnetic flux noise in superconducting devices. Phys. Rev. Applied 6, 044001 (2016).
53. Müller, C., Cole, J. H. & Lisenfeld, J. Towards understanding two-level-systems in amorphous solids: insights from quantum circuits. Rep. Prog. Phys. 82, 124501 (2019).
Acknowledgements
We are grateful to J. Lisenfeld and W. Wulfhekel for insightful discussions as well as A. Lukashenko and L. Radtke for technical assistance. Funding was provided by the Alexander von Humboldt Foundation in the framework of a Sofja Kovalevskaja award endowed by the German Federal Ministry of Education and Research, and by the European Union’s Horizon 2020 programme under no. 899561 (AVaQus). M.S. and G.C. acknowledge support from the German Ministry of Education and Research (BMBF) within the project GEQCOS (FKZ: 13N15683 and 13N15685). P.P. acknowledges support from the German Ministry of Education and Research (BMBF) within the QUANTERA project SiUCs (FKZ: 13N15209). D.R., S.G. and W.W. acknowledge support from the European Research Council advanced grant MoQuOS (no. 741276). Facilities use was supported by the KIT Nanostructure Service Laboratory. We acknowledge qKit for providing a convenient measurement software framework.

Author contributions
M.S. and I.M.P. conceived the experiment. N.D. and Y.C. installed and supervised the real-time microwave electronics. M.S. performed the experiment, analysed the data and developed the theoretical model. I.T. and P.W. designed and fabricated the parametric amplifier. R.G. and O.S. provided the real-time microwave electronics for preliminary experiments. P.P. and N.G. carried out additional measurements for the revised manuscript. M.S. and I.M.P. wrote the manuscript. A.S. and I.M.P. supervised the project. All authors discussed the results and contributed to the final manuscript.

Competing interests
The authors declare no competing interests.

Additional information
Correspondence and requests for materials should be addressed to Martin Spiecker or Ioan M. Pop.