

Probing an electron-spin ensemble with squeezed microwave signals

P. Oehrl ^{1,2} F. Fesquet,^{1,2} K. E. Honasoge ^{1,2} M. Handschuh ^{1,2} A. Marx,¹ R. Gross ^{1,2,3}
K. G. Fedorov ^{1,2,3} and H. Huebl ^{1,2,3,*}

¹*Walther-Meißner-Institut, Bayerische Akademie der Wissenschaften, 85748 Garching, Germany*

²*Department of Physics, TUM School of Natural Sciences, Technical University of Munich, 85748 Garching, Germany*

³*Munich Center for Quantum Science and Technology, 80799 Munich, Germany*



(Received 28 December 2025; accepted 7 July 2026; published 4 August 2026)

The efficient transfer of quantum states into a long-lived storage unit such as solid-state spin ensembles is widely recognized as a critical challenge with significant implications for quantum communication, sensing, and computing applications. Here, we experimentally investigate the interaction of propagating squeezed microwaves with an electron-spin-resonance transition in order to evaluate the use of spin ensembles as quantum memories for gigahertz signals. We generate continuous-variable microwave states with squeezing levels up to 5 dB below the vacuum level and let this signal interrogate a spin ensemble, which is inductively coupled to a lumped-element superconducting microwave resonator with a cooperativity of $C = 0.3$. By analyzing this signal using Wigner tomography and modeling our experimental results with a dedicated steady-state model based on the quantum input-output formalism, we extract a transfer efficiency of around 61 % between the squeezed microwaves and the spin excitations, resulting in squeezing levels of the spin mode up to 0.5 dB below the vacuum level.

DOI: [10.1103/f81b-px11](https://doi.org/10.1103/f81b-px11)

I. INTRODUCTION

The realization of secure communication in future quantum networks requires high-fidelity transfer and storage of quantum states [1–4]. In particular, long-range quantum key distribution and quantum repeater schemes are expected to rely on long-lived quantum memories (QMs) to preserve non-classical or entangled states and offset inevitable temporal delays in practical networks [5–7]. Another promising application of QMs in communication protocols concerns the idea of quantum tokens, where quantum states are associated with unique identities or credentials, which can be securely verified through communication with a trusted node [8–10]. The security of such schemes relies on fundamental information-theoretic assumptions, which provide access to the concept of unconditional security, in contrast to the computational security provided by modern classical encryption schemes.

Among various QM platforms [11–13], solid-state spin ensembles are considered to be very promising candidates due to their long coherence times and ability to be coupled to both optical and microwave signals. In the last decade, significant progress towards the implementation of long-lived QMs with spin ensembles has been achieved by demonstrating long coherence times up to several hours [14,15], the potential

for scalability and for integration into networks [16,17], the ability to store and retrieve coherent states in the limit of single photons [18,19], and the realization of a random-access QM with classical microwave signals [20]. Nonetheless, the realization of efficient quantum state storage and retrieval at microwave frequencies in the single-photon regime remains an open challenge [20–22]. The experimental demonstration of such a microwave QM platform would enable further integration with superconducting quantum circuits, one of the leading architectures in quantum information processing.

Here, we present experimental results based on a modular combination of a Josephson parametric amplifier (JPA) for the generation of squeezed microwave signals and a spin ensemble coupled to a microwave cavity for the QM implementation. We demonstrate the successful generation of propagating microwave squeezed states with squeezing levels up to 5.3 dB below vacuum. We couple these nonclassical photonic states to the excitations of a spin-resonator hybrid and study the resulting response in the steady-state regime. Based on an input-output formalism, we estimate a transfer efficiency of 61% between the quantum microwave states and spin excitations. Relying on our model in the steady-state regime, we can compute a corresponding squeezing of the spin mode of 0.5 dB below vacuum fluctuations.

II. EXPERIMENTAL SETUP

A conceptual scheme of our experimental setup is presented in Fig. 1. We generate squeezed vacuum states at microwave frequencies using a reflection-type, flux-driven JPA consisting of a coplanar waveguide resonator coupled to a nonlinear Josephson-junction circuit [23–25]. This circuit

*Contact author: hans.huebl@wmi.badw.de

Published by the American Physical Society under the terms of the Creative Commons Attribution 4.0 International license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI.

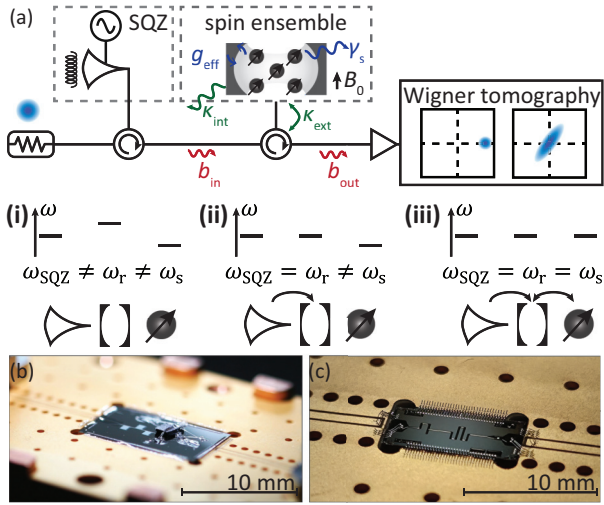


FIG. 1. Experimental setup. (a) By sending a vacuum state to the phase-sensitively operated JPA, a squeezer (SQZ), we generate a propagating squeezed vacuum signal $\hat{b}_{\text{in}}$ with frequency ω_{SQZ} , which is directed to the spin-resonator system via superconducting coaxial cables and a cryogenic circulator. The transfer of the squeezed signal to the spin excitation is mediated via a superconducting lumped-element microwave resonator. The output signal from the spin-cavity system $\hat{b}_{\text{out}}$ propagates further through a cascade of cryogenic and room-temperature amplifiers and is then down-converted and digitized to perform the Wigner tomography of microwave signals. By tuning the static magnetic field B_0 , we can distinguish three coupling cases: (i) The SQZ is off resonant from the spin-resonator hybrid, (ii) the SQZ is in resonance with the resonator but off resonant with the spin ensemble, and (iii) the SQZ, resonator, and spin ensemble are tuned in resonance. (b) Photo of the spin-resonator system. (c) Photo of the JPA device.

is realized by a dc superconducting quantum interference device (SQUID), which acts as a flux-tunable inductor and thus changes the JPA resonance frequency ω_{SQZ} as a function of the applied local magnetic flux. The JPA resonator operates in the overcoupled limit and uses an input circulator which separates the input vacuum state from output squeezed states. In addition, the JPA has an independent microwave port, inductively coupled to the dc SQUID, which is used to apply a strong parametric pump signal at the frequency $2\omega_{\text{SQZ}}$. More details about the design, fabrication, and performance of the JPA can be found in Ref. [26].

The propagating nonclassical squeezed signal is directed via a microwave cryogenic circulator to the spin-resonator hybrid device, which resembles the quantum memory element in our setup. After interacting with the spin-resonator system, the reflected signal is amplified by a broadband low-noise cryogenic high-electron-mobility transistor (HEMT) amplifier. The amplified output signal is processed further at room temperature by using a heterodyne detection setup with a high-frequency digitizer [27]. We calibrate the gain and noise of the amplification chain, including the JPA properties, by making use of the Planck spectroscopy [28,29]. This calibration procedure allows us to perform the Wigner tomography of our quantum states at various points of the cryogenic setup based on the reference-state-reconstruction approach [30].

The hybrid spin-resonator system is composed of a reflection-type planar niobium microwave resonator and the electron-spin ensemble, provided by an isotopically enriched ^{28}Si crystal with a reduced nuclear-spin population that is doped with phosphorus donors at a concentration of $[P] = 1 \times 10^{17} \text{ cm}^{-3}$ [31]. The 18- μm -thick ^{28}Si phosphorus-doped crystal is mounted in a flip-chip configuration onto the lumped-element niobium resonator and held in place using a second silicon chip with a natural isotope composition acting as a cover. While phosphorus donors in silicon exhibit two electron-spin-resonance transitions due to the Fermi-contact hyperfine interaction [32,33], we will selectively exploit only the $|\downarrow\downarrow\rangle \rightarrow |\uparrow\downarrow\rangle$ transition with frequency ω_s . Here, we use the notation $|\downarrow\downarrow\rangle$ to denote the energy levels, where the first (second) arrow corresponds to the electron (nuclear) spin. The energy levels are labeled according to their high-field product-state configuration (see also Appendix A).

By controlling the static magnetic field B_0 , provided by a superconducting solenoid magnet, ω_s can be tuned into resonance with the frequency of the resonator ω_r . Notably, the microwave resonator also shows a weak frequency dependence on the applied in-plane magnetic field, which we attribute to the field dependent kinetic inductance of niobium. The JPA is spatially separated from the spin-resonator hybrid and placed outside the magnet. The additional superconducting magnetic shield guarantees close to quantum-limited performance. This setting allows us to control the resonance frequency of the microwave resonator and the spin ensemble by tuning the external magnetic field B_0 independently of the JPA frequency.

Ultimately, we can distinguish between three different scenarios: (i) a completely off-resonant regime, where the frequencies of the JPA, microwave resonator, and spin ensemble are all different, $\omega_{\text{SQZ}} \neq \omega_r \neq \omega_s$; (ii) a partially resonant regime, where the squeezed signal from the JPA (SQZ) is resonant with the microwave resonator but the spin ensemble remains detuned, $\omega_{\text{SQZ}} = \omega_r \neq \omega_s$; and (iii) a fully resonant regime, where all characteristic carrier frequencies of the aforementioned systems coincide, $\omega_{\text{SQZ}} = \omega_r = \omega_s$. These three distinct configurations allow us to unambiguously characterize all parts of our cryogenic setup and calibrate all important quantities independently, such as the squeezing level of the incoming microwave signals, the microwave response of the resonator, and the spectral characteristics of the spin-resonator hybrid system.

III. THEORY

Here, we briefly describe our approach to model the spin-resonator hybrid system consisting of a driven microwave resonator coupled resonantly to an electron-spin ensemble within a steady-state, quantum-mechanical framework. It enables the analysis of the mean-field amplitudes as well as the variances of microwave quantum signals.

We define the Hamiltonian of the spin-resonator hybrid as

$$\begin{aligned} \hat{H}/\hbar = & \omega_r \hat{a}^\dagger \hat{a} + \omega_s \hat{s}^\dagger \hat{s} + g_{\text{eff}} (\hat{s}^\dagger \hat{a} + \hat{s} \hat{a}^\dagger) \\ & + i\sqrt{2\gamma_s} (\hat{s}^\dagger \hat{f}_s + \text{H.c.}) + i(\sqrt{2\kappa_{\text{int}}} \hat{b}_1 \hat{a}^\dagger + \text{H.c.}) \\ & + i(\sqrt{2\kappa_{\text{ext}}} \hat{b}_{\text{in}} \hat{a}^\dagger + \text{H.c.}), \end{aligned} \quad (1)$$

where we assume weak excitation signals and also neglect the inhomogeneous spin broadening [22,34–36]. Here, $\hat{a}$ ($\hat{s}$) and $\hat{a}^\dagger$ ($\hat{s}^\dagger$) are the bosonic annihilation and creation operators for excitations inside the resonator (spin ensemble), fulfilling $[\hat{a}(t), \hat{a}^\dagger(t)] = 1$ ($[\hat{s}(t), \hat{s}^\dagger(t)] = 1$). The collective coupling between the spin ensemble and the resonator is defined by $g_{\text{eff}} = \sqrt{N}g_0$, with N being the number of spins and g_0 being the single-spin coupling rate. The eigenfrequency of the resonator (spin ensemble) is ω_r (ω_s). The spin dephasing is modeled using a collective coupling with rate γ_s to a noise bath $\hat{f}_s$. Similarly, we account for internal resonator losses with an associated loss mode $\hat{b}_1$ and a loss rate κ_{int} . The coupling between the intracavity mode of the resonator and the external microwave circuit is described via the coupling of the input mode $\hat{b}_{\text{in}}$ with rate κ_{ext} . We assume all coupling rates are half width at half maximum linewidth.

Using the Heisenberg equation of motion, we derive the following equations for the Fourier amplitudes of the resonator and spin operators:

$$-i\omega\hat{a}(\omega) = -(i\omega_r + \kappa)\hat{a}(\omega) + \sqrt{2\kappa_{\text{int}}}\hat{b}_1(\omega) + \sqrt{2\kappa_{\text{ext}}}\hat{b}_{\text{in}}(\omega) - ig_{\text{eff}}\hat{s}(\omega), \quad (2)$$

$$-i\omega\hat{s}(\omega) = -(i\omega_s + \gamma_s)\hat{s}(\omega) - ig_{\text{eff}}\hat{a}(\omega) + \sqrt{2\gamma_s}\hat{f}_s(\omega), \quad (3)$$

with $\kappa = \kappa_{\text{int}} + \kappa_{\text{ext}}$. Following Ref. [22], we obtain the relation between the input and output signal variances as

$$\sigma_{\hat{b}_{\text{out}}}^2(\omega) = |r(\omega)|^2\sigma_{\hat{b}_{\text{in}}}^2 + |l(\omega)|^2\sigma_{\hat{b}_1}^2 + |t(\omega)|^2\sigma_{\hat{f}_s}^2. \quad (4)$$

Here, the variance of an operator $\hat{O}$ is defined as $\sigma_{\hat{O}}^2 = \langle \hat{q}_{\hat{O}}^2 \rangle$, with $\hat{q}_{\hat{O}} = (\hat{O} + \hat{O}^\dagger)/2$. The complex coefficients r , l , and t describe terms of the Langevin noise operator associated with the reflection from the cavity, the resonator internal losses, and losses of the spin ensemble, respectively. These coefficients are [22]

$$\begin{aligned} r(\omega) &= \text{Re}\left(1 - \frac{2\kappa_{\text{ext}}}{X(\omega)}\right), \\ l(\omega) &= \text{Re}\left(\frac{\sqrt{4\kappa_{\text{ext}}\kappa_{\text{int}}}}{X(\omega)}\right), \\ t(\omega) &= \text{Re}\left(\frac{Y(\omega)}{X(\omega)}\right), \end{aligned} \quad (5)$$

with the following definitions:

$$\begin{aligned} X(\omega) &= -i\omega + \kappa + \frac{g_{\text{eff}}^2}{[\gamma_s + i(\Delta_{\text{sr}} - \omega)]}, \\ Y(\omega) &= \frac{\sqrt{4\gamma_s\kappa_{\text{ext}}g_{\text{eff}}}}{\gamma_s + i(\Delta_{\text{sr}} - \omega)}. \end{aligned} \quad (6)$$

Here, we define the detuning of the spin and resonator modes as $\Delta_{\text{sr}} = \omega_s - \omega_r$ and the detuning of the resonator mode and the incident microwave signal as $\Delta_r = \omega - \omega_r$. The condition $|l(\omega)|^2 + |r(\omega)|^2 + |t(\omega)|^2 = 1$ ensures that the output mode fulfills the bosonic commutation relation. We note that Eq. (4) with the defined coefficients r , l , and t is valid only at zero or large spin-resonator detunings [37]. In addition, we compare the expectation value of the input and output fields within the input-output formalism and derive the complex scattering

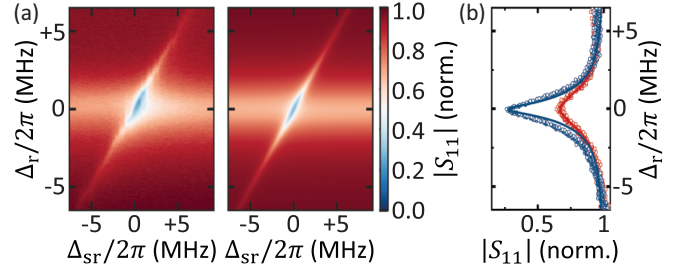


FIG. 2. Characterization of the spin-resonator hybrid. (a) Measured (left) and modeled (right) microwave reflection amplitudes $|S_{11}|$ as a function of the coherent probe-resonator detuning Δ_r and spin-resonator detuning Δ_{sr} . The spectrum is modeled based on Eq. (7). (b) Comparison of the measured (dots) and modeled (lines) data for the resonant ($\Delta_{\text{sr}}/2\pi = 0$) and off-resonant ($\Delta_{\text{sr}}/2\pi = 9.2$ MHz) conditions, presented in blue and red, respectively.

parameter [38–40]

$$S_{11}(\omega) = \frac{\langle \hat{b}_{\text{out}} \rangle}{\langle \hat{b}_{\text{in}} \rangle} = 1 + \frac{2\kappa_{\text{ext}}}{i(\omega_r - \omega) - \kappa + \frac{g_{\text{eff}}^2}{i(\omega_s - \omega) - \gamma_s}}. \quad (7)$$

We analyze the output-field variance in order to determine the squeezing level S of the corresponding quantum state. To this end, we describe the output microwave signal in terms of its electromagnetic field quadrature operators, $\hat{q}$ and $\hat{p}$, with $\hat{b}_{\text{out}} = \hat{q} + i\hat{p}$, fulfilling the commutator relation $[\hat{q}, i\hat{p}] = 1/2$. We can completely define the squeezed and antisqueezed quadrature variances within this nomenclature, as derived in Ref. [30]. Consequently, the squeezing level can be expressed in decibels as

$$S = -10 \log_{10}(4\sigma_{\text{sq}}), \quad (8)$$

where σ_{sq} is the variance of the squeezed quadrature and positive values indicate squeezing below the vacuum level [26].

IV. CHARACTERIZATION OF SPIN-RESONATOR HYBRID

Next, we perform continuous-wave microwave spectroscopy of the resonator coupled to the electron-spin-resonance (ESR) transition. In detail, we measure the complex scattering parameter S_{11} as a function of frequency using an incident microwave coherent signal with a power of $P_{\text{mw}} = -139$ dBm, corresponding to an average resonator occupation of 0.5 photon at ω_r [41]. By controlling the applied static magnetic field B_0 , we can tune the frequency of the ESR transition to match the resonator resonance frequency, $\omega_s = \omega_r$. For this condition, S_{11} also reflects the conversion of coherent microwave photons to spin excitations. Figure 2(a) shows the measured microwave reflection amplitude $|S_{11}|$ close to $\omega_r \simeq \omega_s$ as a function of the detunings $\Delta_{\text{sr}} = (\omega_s - \omega_r) \propto B_0 - B_{\text{res}}$ and $\Delta_r = (\omega - \omega_r)$. In the two-dimensional plot we identify two distinct features: One is an absorption line at $\Delta_r \approx 0$ which is independent of Δ_{sr} , i.e., independent of the applied magnetic field, and is associated with the fixed-frequency microwave resonator; we also observe an ESR absorption signature which varies linearly with the applied magnetic field, i.e., Δ_{sr} . We can

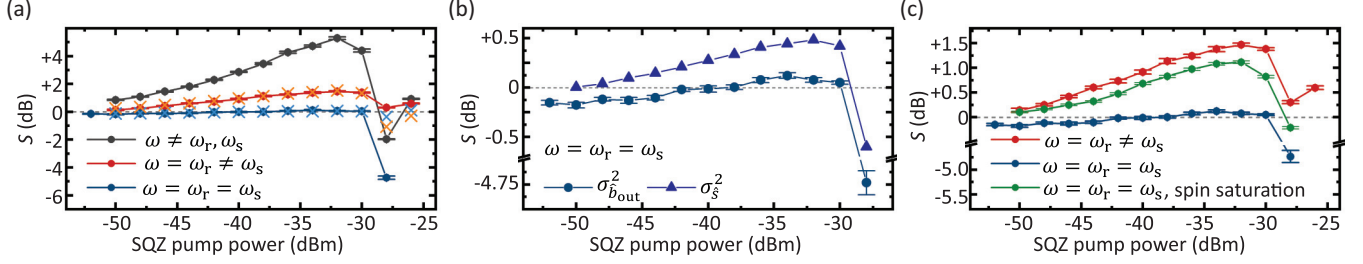


FIG. 3. Power-dependent squeezing levels S of propagating microwave signals and the spin mode. (a) The measured squeezing level S as a function of the SQZ pump power in the off-resonant case (black circles), in resonance with the resonator but off resonant with the spins (red circles), and in resonance with the spin-resonator hybrid (blue circles). The lines are guides to the eyes. The orange and light-blue crosses represent our model predictions based on Eq. (4) using the off-resonant squeezing states as reference input signals combined with independently determined coupling parameters g_{eff} , κ_{ext} , κ_{int} , and γ_s . (b) For the resonant case, the measured squeezing level of the output microwave signals $\sigma_{\text{b\_out}}^2$ indicates squeezing below vacuum and is significantly reduced due to a transfer of the microwave states to spin excitations. Based on an effective rate model, the variance of the spin mode σ_s^2 can be calculated, resulting in squeezing levels up to 0.5 dB. (c) Spin saturation measurements (green) using a strong resonant microwave pulse that saturates the ESR transition before performing a continuous-wave squeezing measurement. This results in an increase of the squeezing levels approaching the reference experiment.

determine the characteristic parameters of both subsystems by analyzing these spectral features. For $\Delta_{\text{sr}}/2\pi = 9.2$ MHz, which is far detuned from the ESR line, we extract the resonance frequency $\omega_r/2\pi = 5.645$ GHz, as well as the resonator linewidth $\kappa/2\pi = 1.78$ MHz and external coupling constant $\kappa_{\text{ext}}/2\pi = 1.49$ MHz, using the circle-fit method (see Appendix C) [42]. We next fit the total measured spectrum using Eq. (7), treating the aforementioned resonator parameters as known input parameters. In this way, we determine the intrinsic spin relaxation rate to $\gamma_s/2\pi = 382$ kHz and the effective coupling rate to $g_{\text{eff}} = 460$ kHz. The experimentally determined cooperativity is $C = g_{\text{eff}}^2/(\gamma_s\kappa) = 0.3$.

We compare the continuous-wave ESR (cwESR) spectroscopy data with the response spectrum predicted by the input-output model in Eq. (7), which we present in Fig. 2(a). In addition, Fig. 2(b) presents selected cross sections for the resonant and off-resonant cases of the spectra, demonstrating quantitative agreement between the model prediction and the experimental data. We note a slight asymmetry in the experimentally measured absorption dips which is not described by our model and which is not corrected by the subtraction of the complex microwave background induced from imperfections of the experimental setup (see also Appendix C). Moreover, using conventional pulsed electron-spin-resonance techniques, we determine the coherence time $T_2 = 2.16(0.11)$ ms and the spin relaxation time $T_1 = 85.49(0.99)$ s (see Appendix A).

V. PROBING THE SPIN-RESONATOR HYBRID SYSTEM WITH SQUEEZED LIGHT

Next, we experimentally investigate the reflection of propagating microwave squeezed states from the spin-resonator hybrid. To this end, we generate squeezed vacuum states using the squeezer JPA (SQZ), guide these non-classical microwaves towards the input circulator of the spin-resonator system, and analyze the reflected signals using the reference-state-reconstruction protocol [30]. We measure the amplitudes of amplified microwave output signals using a field-programmable gate-array (FPGA) module with a

detection bandwidth of 400 kHz. We use the FPGA internal logic to perform averaging and calculate the moments of the microwave signals up to fourth order. This approach allows us to derive the Wigner function of the involved quantum states in the Gaussian approximation and estimate the squeezing level S as a function of the applied pump power [30,43].

Figure 3(a) shows the measured squeezing levels for the three different coupling scenarios: (i) $\omega_{\text{SQZ}} \neq \omega_r \neq \omega_s$, (ii) $\omega_{\text{SQZ}} = \omega_r \neq \omega_s$, and (iii) $\omega_{\text{SQZ}} = \omega_r = \omega_s$. Note that we experimentally determine the squeezing level at the input of the HEMT amplifier (see Appendix C for details). All measurements are performed for a fixed frequency $\omega_{\text{SQZ}}/2\pi = 5.645$ GHz. At $B_0 = 0$ the weak field dependence of the resonator results in a detuning of $(\omega_r - \omega_{\text{SQZ}})/2\pi \simeq 10$ MHz [44], compared to the spin-ensemble working point at $\Delta_{\text{sr}} = 0$. This results in a full reflection of the squeezed microwave state from the resonator input and allows us to obtain a reference baseline for the squeezing level as a function of the pump power. In this configuration, we observe a steady increase of S with increasing pump power up to a maximum squeezing level of $S_{\text{(i), max}} = (+5.29 \pm 0.09)$ dB. Here, positive squeezing values indicate squeezing below the vacuum level. Above a pump power of -30 dBm, the squeezing level abruptly decreases and eventually crosses the threshold value of $S = 0$ dB. This could originate from the reconstruction model considering moments up to fourth order or from the disappearance of nonclassical correlations in the signal. In the next step, we increase the static magnetic field, such that $\Delta_{\text{sr}}/2\pi = +9.2$ MHz. Then, the frequencies of the SQZ and the resonator coincide while the spin mode remains detuned. In this configuration, the pump-power dependence of S is qualitatively similar to the off-resonant reference measurement but shows a reduction in the squeezing level. The measured maximum squeezing level is $S_{\text{(ii), max}} = (+1.47 \pm 0.03)$ dB. This reduction originates from the intrinsic resonator losses, as determined in the cwESR spectroscopy (see Fig. 2). These losses lead to an effective addition of thermal fluctuations from the coupled environmental bath, and consequently lead to an eventual reduction of the squeezing level. Finally, we tune all our subsystems in resonance,

$\omega_{\text{SQZ}} = \omega_r = \omega_s$, and measure the resulting squeezing level. We observe a further reduction of the power-dependent squeezing level as compared to the previous configurations. Here, the measured maximum squeezing level is $S_{\text{(iii)}, \text{max}} = (+0.12 \pm 0.03)$ dB. We attribute this additional reduction of the squeezing level to the increased total absorption losses associated with the ESR transition.

For a quantitative analysis, we compare the presented experimental results with the prediction of the input-output model, as given by Eq. (4). Within this model, the variances in the measured output signal are determined by the variances of the reflected incoming squeezed microwave signals and thermal-state contributions from the resonator and spin-ensemble modes, which are assumed to be weak thermal baths. As an input signal variance $\sigma_{b_{\text{in}}}$, we use the value obtained from the reference measurement, where $\omega_{\text{SQZ}} \neq \omega_r \neq \omega_s$. Using the coupling rates obtained from the earlier cwESR measurements and describing the spin and resonator losses as arising from a thermal bath with the same temperature $T = 70$ mK, we can use Eq. (4) to estimate the expected squeezing level from the measured output variance. The calculated values are presented as crosses in Fig. 3(a), showing excellent quantitative agreement with the measured squeezing levels. Furthermore, the parameter t in Eq. (4) can be interpreted as a transfer efficiency between microwave signals and spin excitations, as it describes how strong thermal fluctuations of the spin systems couple to the microwave signal. Based on the fit of our model to the experimental data, we find $|t(0)|^2 \simeq 61\%$.

Within the steady-state regime, we can also derive an effective rate model to evaluate the quadrature variances of our spin ensemble (see Appendix E). Within this theoretical description we have to assume that the spin system can be modeled as a mode with a finite bandwidth of $\gamma_s/2\pi = 382$ kHz. Focusing on the steady-state level of fluctuation (i.e., we do not capture transient, time-dependent dynamics), we assume that the spin ensemble is coupled to an effective bath that can be described by a thermal photon number and a squeezing parameter [45–47]. As a result, we can compute the variances of the collective spin excitations and determine the maximum squeezing level of the spin mode to be 0.5 dB below vacuum for the input squeezing of +5.29 dB. Figure 3(b) shows the power-dependent spin-squeezing levels. In combination with our measurements of the signal output variance, our model suggests that we successfully transferred a nonclassical microwave state to an excitation of a spin ensemble. Further confirmation requires time-domain investigations, e.g., by combining pulsed techniques for nonclassical microwave signal generation with quantum memory protocols [20,22].

Complementary to detuning the subsystems, another method to decouple the spin ensemble from the microwave resonator exists. Conceptually, the application of a strong resonant microwave pulse can saturate the spin transition and decouple the resonator from the spin ensemble. In order to experimentally investigate this regime, we operate in the resonant configuration between all three subsystems and additionally apply a resonant, strong, rectangular-shaped, saturation pulse at a carrier frequency of $\omega/2\pi = 5.645$ GHz for a duration of 20 ms with a power of -120 dBm

before performing the continuous-wave measurements with the squeezed states. Within our squeezing measurements, we choose a pulse repetition time of 2 s. The measured squeezing level at the output of the spin-resonator hybrid is shown in Fig. 3(c), with a maximum squeezing level of $S_{\text{spins, sat}} = (+1.12 \pm 0.03)$ dB. This value reaches close to the originally measured maximal squeezing level for the case of the SQZ in resonance with the resonator without the saturation pulse. The clear increase in squeezing level thus confirms that the saturation pulse leads to a decoupling of the spin ensemble and the microwave resonator. Quantitatively, the squeezing levels do not reach the reference values (red line in Fig. 3). We attribute this observation to two aspects: the nonlinear properties of the spin-resonator hybrid and the timescales of the saturation protocol, i.e., the finite T_1 time of the donor spin ensemble which counteracts the spin-resonator decoupling. These nonlinear properties require a more complex description going beyond the simple model provided by Eq. (1).

VI. SIMULATION OF THE QUANTUM TRANSFER EFFICIENCY

In order to investigate optimal coupling parameters for the spin-resonator hybrid platform, we numerically compute the transfer efficiency using Eq. (4). We study experimentally reachable parameter regimes for realistic internal and external resonator coupling strengths, the spin dephasing rate, and an effective spin-resonator coupling. Figure 4 presents the results of those numerical calculations, which provide a foundation for the realization of an optimal quantum state transfer for future experiments. Here, we investigate the resonant case, $\omega_{\text{SQZ}} = \omega_r = \omega_s$, and thus investigate the transfer efficiency $|t(0)|^2$ as a function of different coupling rates.

In the first case, we choose the spin dephasing rate to be fixed, $\gamma_s = 382$ kHz, and vary the external coupling rate of the resonator as well as the effective spin-resonator coupling. Figures 4(a) and 4(b) present the estimated effective transfer strength and cooperativity corresponding to the maximized transfer efficiency, respectively. We observe a parameter regime with increasing coupling efficiency, approaching unity in the case of $\kappa_{\text{ext}}/2\pi = 1.25$ MHz and $g_{\text{eff}}/2\pi = 976$ kHz. For these parameters, the corresponding cooperativity reaches $C_{\text{max}} = 0.982$. Next, we fix the spin-resonator coupling, $g_{\text{eff}} = 460$ kHz, and vary the external coupling of the resonator as well as the spin dephasing rate. Figures 4(c) and 4(d) show the resulting coupling regime and corresponding cooperativity. The coupling efficiency reaches a value of 0.991 for $\kappa_{\text{ext}}/2\pi = 1.25$ MHz and $\gamma_s/2\pi = 167$ kHz, almost corresponding to unit cooperativity. In the third scenario, we keep the spin dephasing rate and effective spin-resonator coupling constant and equal to the determined experimental values, i.e., $\gamma_s = 382$ kHz and $g_{\text{eff}} = 460$ kHz. Then, we vary the internal and external coupling rates of the resonator, which can experimentally be achieved by altering the resonator design. We find the maximum coupling efficiency for vanishing internal resonator losses ($\kappa_{\text{int}} = 0$ kHz) and $\kappa_{\text{ext}}/2\pi = 556$ kHz with a corresponding cooperativity of $C = 0.999$.

Our simulations of the coupling efficiency for various coupling parameters of our spin-resonator hybrid system are in

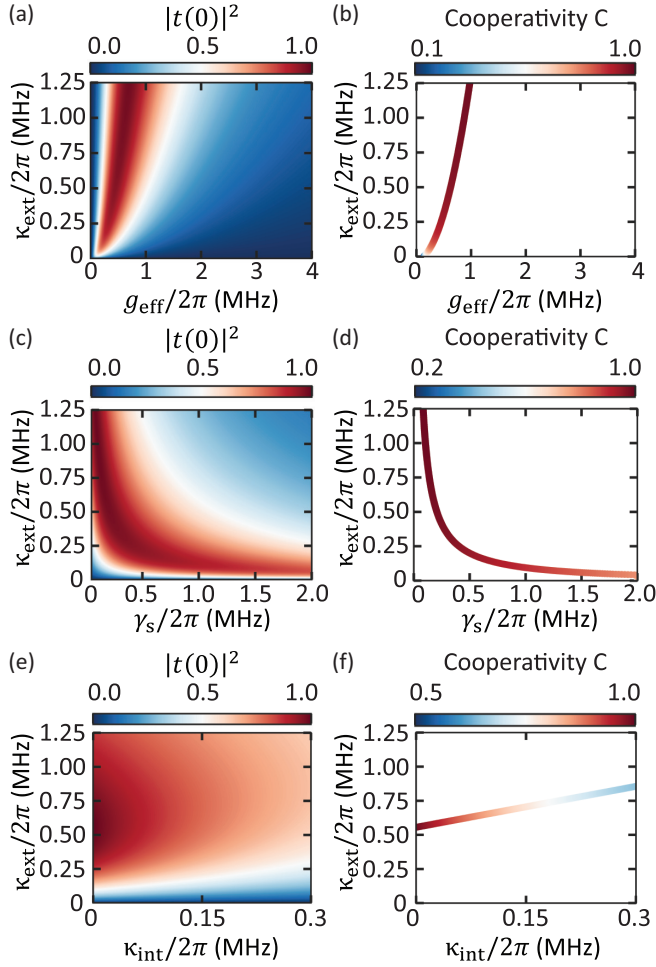


FIG. 4. Numerical calculation of the quantum transfer efficiency $|t(0)|^2$. The transfer efficiency is computed according to Eq. (4) as a function of different coupling strengths for the resonance case, i.e., $\omega_{\text{SQZ}} = \omega_r = \omega_s$. (a) Transfer efficiency as a function of the external coupling rate κ_{ext} and effective spin-resonator coupling g_{eff} for $\gamma_s/2\pi = 382$ kHz. (c) Transfer efficiency as a function of the external coupling rate κ_{ext} and spin dephasing γ_s for $g_{\text{eff}}/2\pi = 460$ kHz. The spin dephasing rate and the spin-resonator coupling are fixed to $\gamma_s = 382$ kHz and $g_{\text{eff}} = 460$ kHz, respectively. (e) Transfer efficiency as a function of the external and internal coupling rates, κ_{ext} and κ_{int} , respectively. Optimal state transfer and maximum transfer efficiency are achieved in the limit of unit cooperativity, $C = 1$. (b), (d), and (f) illustrate the calculated cooperativity for the maximum coupling efficiency.

agreement with results reported in the literature, as can be found in Refs. [48–50]. We want to highlight that our spin-resonator hybrid system is engineered with a cooperativity of $C = 0.3$, for which we expect a high coupling efficiency. In order to further increase the coupling efficiency of quantum states to our spin ensemble, we can tune the coupling parameters of our microwave resonator. One way can be to improve the internal quality factor, e.g., by using magnetic-field-robust superconducting materials like NbTiN [44]. Furthermore, we can adapt the design of the microwave circuit to tune the coupling of incoming quantum states to the microwave

resonator, e.g., by changing the distance between the hanger-type resonator and the signal feed line.

VII. CONCLUSION

In summary, we studied a spin-resonator hybrid system interacting with squeezed microwave signals at various conditions. We observed a strong transfer of the squeezed microwaves to the spin ensemble corresponding to an efficiency of $|t|^2 \simeq 61\%$. We successfully modeled our experimental findings using an input-output quantum formalism which also allows us to extrapolate the resonator properties required to reach the cooperativity of unity, as demanded by the optimal state transfer with spin ensembles [50,51]. While remaining in the steady-state regime, we can use an effective reservoir description and determine squeezing levels of the spin mode up to +0.5 dB below vacuum. This spin squeezing is induced by the coupling of propagating squeezed microwave signals to our spin-resonator hybrid system. Thus, we demonstrated an actively controllable way for the generation of squeezing in the mode of the spin ensemble.

Our results provide insight into the fundamental understanding of cavity-based ESR with solid-state spin systems driven with nonclassical signals. The evaluation of input- and output-field quadratures in combination with the steady-state solution of the input-output formalism highlights both an important step towards on-demand spin squeezing for quantum sensing applications and the realization of quantum memories for microwave quantum communication and general quantum information processing [52]. The natural next step is the experimental investigation and retrieval of the state stored in the spin excitation, e.g., using time-domain detection concepts [20].

ACKNOWLEDGMENTS

We acknowledge support from the German Research Foundation via Germany's Excellence Strategy (Grant No. EXC-2111-390814868), the German Federal Ministry of Education and Research via the project QuaMToMe (Grant No. 16KISQ036), and the German Research Foundation via the Transregio ConQuMat TRR 360 (Project ID No. 492547816). This research is part of the Munich Quantum Valley (lighthouse project IQ-Sense), which is supported by the Bavarian state government with funds from the Hightech Agenda Bayern Plus. We are grateful to Syeda Aliya Batool for her support and constructive discussions regarding the theoretical aspects of this work.

DATA AVAILABILITY

The data that support the findings of this article are openly available [53].

APPENDIX A: SPIN PROPERTIES

1. Phosphorus donors in silicon

A single phosphorus donor in silicon provides an unbound electron with electron spin $S = 1/2$ and a nuclear spin from the ^{31}P nuclei of $I = 1/2$. Consequently, a corresponding spin

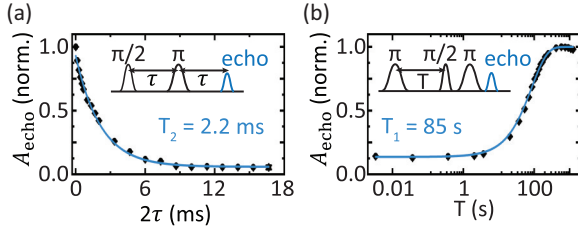


FIG. 5. Coherence times of the Si:P ensemble. Experimentally determined coherence times (a) T_2 and (b) T_1 based on conventional Hahn-echo and inversion-recovery pulse sequences, respectively. The solid lines represent fits according to an exponential function.

Hamiltonian of the spin ensemble in an external magnetic field can be expressed as [54]

$$\hat{H}_s = g_e \mu_B \hat{B}_0 \hat{S} + g_n \mu_n \hat{B}_0 \hat{I} + A \hat{S} \hat{I}, \quad (\text{A1})$$

where the first two terms are the Zeeman interaction of the electron- and nuclear-spin systems, respectively. The last term represents the Fermi-contact hyperfine interaction with the hyperfine constant $A/h = 117.53$ MHz. The g factors of the electron and nuclear spins are $g_e = 1.9985$ and $g_n = 2.2632$, respectively. Furthermore, μ_B and μ_n are the Bohr and nuclear magnetons, respectively.

In the limit of high magnetic fields ($g_e \mu_B B_0 \gg A$), we expect two spin resonant transitions. In addition to these desired ESR transitions, signatures of dangling bond defects located at the Si/SiO₂ interface are visible in the spectrum (see, e.g., [55,56]). The electron-spin subensemble linked to the nuclear spin-down configuration is known to avoid the overlap of multiple ESR transitions and is thus investigated here.

2. Pulsed ESR spectroscopy

We perform pulsed ESR spectroscopy based on Hahn-echo sequences to determine the characteristic coherence times of our spin ensemble [54,57,58]. We rely on the Zeeman shift of the spin mode in resonance with the coupled cavity mode. Then, we generate in-phase and out-of-phase signals of Gaussian-shaped pulses using an arbitrary waveform generator and modulate carrier signals at the pulses to the resonance frequency of our spin-resonator hybrid using a vector signal generator. The detection of the resulting spin echo is performed using a heterodyne down-conversion setup. We analyze the area of the spin echo A_{echo} by fitting a Gaussian function to the data and integrating the area within a time window around the maximum amplitude of the echo given by 3 times the standard deviation.

Furthermore, we optimize the Gaussian input pulses by using detected echo areas. We choose the widths of the $\pi/2$ pulse and π pulse to be equal but the pulse amplitudes to be different. We find a maximum spin-echo area by choosing a pulse width of $2 \mu\text{s}$ and an amplitude ratio of $0.625/1$ between the $\pi/2$ pulse and π pulse. Using those optimized pulse parameters, we vary the delay time τ between the two incoming pulses to determine the characteristic coherence time T_2 within a conventional Hahn-echo pulse sequence, $\pi/2 - \tau - \pi - \tau - \text{echo}$. Figure 5(a) shows the resulting echo area, which is exponentially decreasing for increasing delay times. We fit an

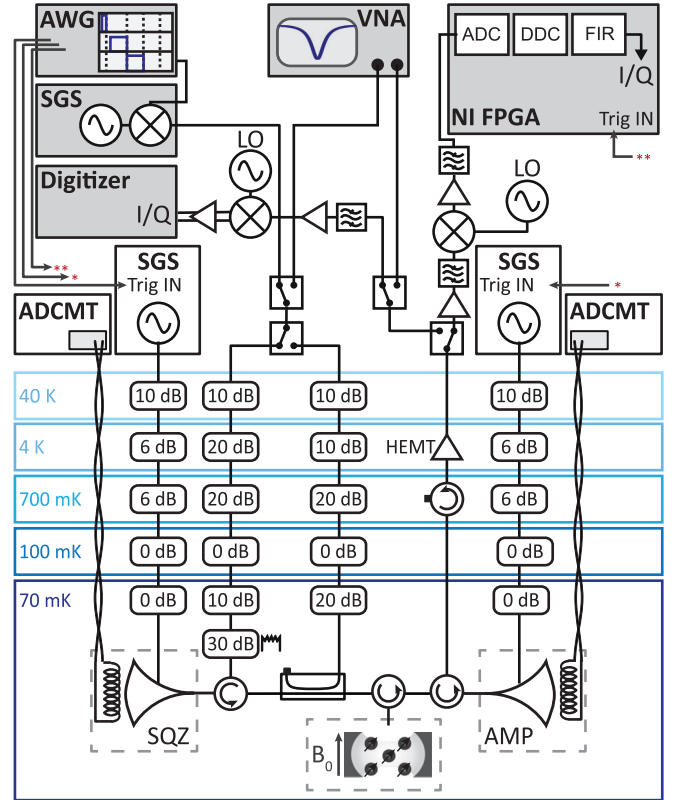


FIG. 6. Experimental setup. Cryogenic and room-temperature setup for continuous-wave and pulsed measurements with coherent microwave signals as well as the generation and detection of squeezed microwave signals. Further details are given in the text.

exponential decay function to the experimental data, resulting in $T_2 = 2.16 \pm 0.11$ ms. Similarly, we perform an inversion recovery pulse sequence, $\pi - T - \pi/2 - \tau - \pi - \tau - \text{echo}$, to determine the characteristic coherence time $T_1 = 85.49 \pm 0.99$ s.

In order to experimentally demonstrate the retrieval of the stored state from the spin ensemble, we can rely on coherent control and pulsed spin-echo protocols for the storage and retrieval of nonclassical states. For a practical implementation, the spin dephasing time T_2 plays a crucial role as it governs the rate at which phase coherence within the spin ensemble decays, thereby limiting the effective duration over which coherent control operations remain accurate. In comparison to Ref. [20], we demonstrate equal T_2 time and a cooperativity that is more than 4 times larger than the reported value. Consequently, we would expect a one-way efficiency 3 times larger than reported in Ref. [20].

APPENDIX B: EXPERIMENTAL SETUP

1. Cryogenic wiring

The experimental setup is shown in Fig. 6. Within the cryogenic microwave circuit, we have a total of four input lines with different attenuations to suppress room-temperature noise. Two of the input lines serve as pump lines for the JPAs. To probe our measurement devices, we have an input line with a total of 90-dB attenuation that is coupled to the SQZ via a cryogenic circulator (Low Noise Factory LNF-CIC4_8A).

We can bypass the SQZ and probe the spin ensemble via an additional input line with 60-dB attenuation, a cryogenic directional coupler (Sirius Microwave CPL2000-18000-30-C), and an additional cryogenic circulator (Low Noise Factory LNF-CIC4_8A).

The output signal is guided through a cryogenic amplification chain via another circulator (Low Noise Factory LNF-CIC4_12A) to a second JPA (AMP) and, ultimately, via a circulator (Quinstar CTH0408KCS), to the cryogenic HEMT amplifier (Low Noise Factory LNF-LNC4_8C) at the 4 K stage. Note that for the experiments presented in this work we did not use the second JPA (AMP) for the amplification of the output signals. As described before, the spin-resonator device is placed in the center field of a solenoid magnet, while both JPAs are placed outside the magnet.

2. Room-temperature setup

At room temperature, we can use a total of three different measurement techniques. For continuous-wave measurements, we probe our devices with a vector network analyzer (Rohde & Schwarz ZVA8). For pulsed ESR measurements, we use an arbitrary waveform generator (Zurich Instruments HDAWG8) to generate in-phase and out-of-phase signals with Gaussian-shaped pulses. The pulses are up-converted to the resonance frequency using a vector signal generator (Rhode & Schwarz SGS100A). For the detection of spin echoes, the output signal is filtered, amplified, and down-converted using an IQ mixer (Marki IQ-0307L). Then, we use an additional variable amplifier (FEMTO DHPVA-200) and digitize the pulsed signal with an analog-to-digital converter (Spectrum M4i.4451-x8). For the generation of squeezed microwave signals, we use a dc voltage source (ADCMT 6241A) to apply a flux bias setting the resonance frequency of the JPA and a microwave signal generator (Rohde & Schwarz SGS100A) for the generation of the pump tone. The detection of the output signal is performed in a heterodyne detection setup, including a variety of bandpass filters, room-temperature amplifiers, and the FPGA module (National Instruments FPGA 7972).

APPENDIX C: DATA PROCESSING AND CALIBRATION TECHNIQUES

1. Background correction of the cwESR spectrum

When cwESR spectra are measured, the measured scattering parameter is subject to the complex microwave background based on imperfections in the experimental setup, such as impedance mismatches and signal delays due to finite cable lengths. To correct for those effects, we apply the circle-fit routine and fit the resonator response detuned from the spin mode, $\Delta_{\text{sr}}/2\pi = 9.2$ MHz, using the following equation [42]:

$$|S_{11, \text{exp}}| = \left| A_{\text{MW}} \left(1 - \frac{2Q_1/Q_{\text{ext}}}{1 + 2iQ_1(\omega/\omega_r - 1)} \right) \right|, \quad (\text{C1})$$

where the complex amplitude $A_{\text{MW}} = a_{\text{mw}} e^{i(\alpha - \omega\tau)}$ captures environmental influences, including amplitude variations, impedance-induced phase shifts, and microwave signal propagation delays. The loaded quality factor Q_1 is defined as the sum of reciprocal values of the internal and external quality factors (Q_{int} and Q_{ext}), $Q_1^{-1} = Q_{\text{int}}^{-1} + \text{Re}(Q_{\text{ext}}^{-1})$. The internal

quality factor accounts for intrinsic losses such as radiative losses, surface resistance, magnetic-field-induced dissipation, and spin-ensemble interactions. The complex external quality factor $Q_{\text{ext}} = |Q_{\text{ext}}| e^{-i\phi}$ quantifies energy dissipation into the external feed line, while the phase ϕ accounts for impedance mismatches and parasitic effects [42]. Applying the circle-fit routine, we find $a_{\text{mw}} = 0.021$, $\alpha = 1.585$, $\tau = 97.35$ ns, $\omega_r/2\pi = 5.645$ GHz, $\phi = -0.002$, $Q_{\text{ext}} = (1900 \pm 11)$, and $Q_1 = (1586 \pm 15)$.

Assuming that the environmental influences are independent of the applied magnetic field for the investigated field range, we correct our measured data as follows:

$$|S_{11, \text{corr}}| = \left| \frac{\text{Re}(S_{11, \text{meas}}) + i\text{Im}(S_{11, \text{meas}})}{a_{\text{mw}} e^{i(\alpha - \omega\tau)}} \right|. \quad (\text{C2})$$

Furthermore, we use a normalization of the corrected data in order to compare the measured spectrum to the theoretical model. The resulting scattering parameter is defined as

$$|S_{11}| = |S_{11, \text{corr}}| / \max(|S_{11, \text{corr}}|). \quad (\text{C3})$$

2. Reference-state reconstruction

As described in the main text, we digitize the microwave output signal at room temperature with a heterodyne detection setup using an FPGA module. With the recorded I/Q sampled data points, we compute the quadrature moments up to fourth order $\langle I^n Q^m \rangle$ for $n + m \leq 4$ [$(n, m) \in \mathbb{N}^2$], and consequently, we compute the corresponding signal moments with respect to a specific reconstruction point in our setup. Thus, we can reconstruct the squeezing angle and squeezing level for the detected output signal [30,59]. In order to map the detected voltage amplitude of the microwave output signal to a photon number, we rely on Planck spectroscopy [28,29]. Here, we use a heatable input attenuator that is placed within our microwave circuit before the SQZ. Then, we measure the output voltage dependent on the temperature of the heatable attenuator. Based on Planck's distribution fitting routine, we can determine a photon-number conversion factor which relates the voltage measured at room temperature to the photon number at the desired reference point in the cryogenic setup. Within our data analysis, we choose the HEMT input as the reconstruction point. To this end, we carefully estimate the microwave losses between the heatable attenuator and the HEMT input to be -2.8 dB.

APPENDIX D: CHARACTERIZATION OF THE SQUEEZER JPA

1. Flux-dependent resonance frequency and nondegenerate gain

Here, we present characterization measurements of the squeezer JPA (SQZ). First of all, we record the dependence of the resonance frequency on the applied magnetic flux. The result is shown in Fig. 7(a). In our experiment, we measure the complex microwave reflection amplitude $|S_{11}|$ as a function of the probe frequency and dc current through the superconducting magnetic field coil. We observe a maximum resonance frequency of $\omega/2\pi = 5.73$ GHz and a tunability range of about 500 MHz. Importantly, we identify two coil currents at which we can tune the SQZ in resonance with our spin-resonator hybrid, fulfilling $\omega_{\text{SQZ}} = \omega_r = \omega_s$. Furthermore, we

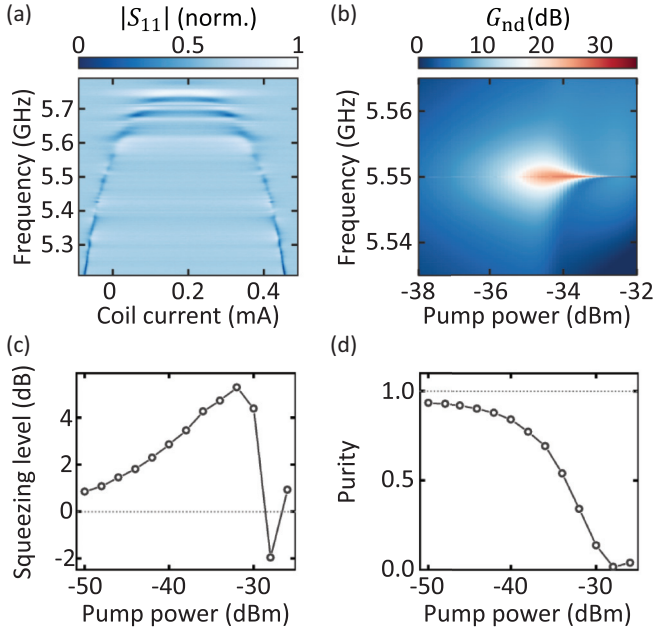


FIG. 7. Experimental characterization of the JPA. (a) Measurement of the microwave reflection amplitude $|S_{11}|$ as a function of the probe frequency and dc current through the superconducting magnetic field coil. (b) Measurement of the dependence of the nondegenerate gain G_{nd} on the applied pump power for a fixed frequency of $\omega_{\text{pump}}/2\pi = 11.1$ GHz. (c) Measured squeezing level as a function of the applied pump power for a fixed frequency of $\omega_{\text{SQZ}}/2\pi = 5.645$ GHz showing a maximum squeezing level of $(+5.29 \pm 0.09)$ dB. The gray dashed line depicts the vacuum limit, above which the reconstructed microwave signals are referred to as squeezed. (d) The corresponding reconstructed purity dependent on the applied pump power. The gray dashed line illustrates the upper limit at which the purity becomes unity.

measure the dependence of the nondegenerate gain G_{nd} on the applied pump power for a fixed frequency of $\omega_{\text{pump}}/2\pi = 2\omega_{\text{SQZ}}/2\pi = 11.1$ GHz. The result is shown in Fig. 7(b). We record a maximum nondegenerate gain of $G_{nd} = 36$ dB for a fixed power of the coherent pump tone of -32.3 dBm.

2. Generation of squeezed microwave signals

Finally, we characterize the generation of squeezed microwave signals by recoding the dependence of the squeezing level on the applied pump power for a fixed frequency of $\omega_{\text{SQZ}}/2\pi = 5.645$ GHz. The results are shown in Fig. 7(c). We observe a maximum squeezing level of $(+5.29 \pm 0.09)$ dB. Correspondingly, we reconstruct the purity as a function of the applied pump power. The results are shown in Fig. 7(d). With increasing pump power, the purity steadily decreases due to pump-induced noise, higher-order nonlinearities, and gain-dependent environmental noise [25,60]. We use this power-dependent squeezing measurement as a reference for the coupling of squeezed microwave signals to the spin-resonator hybrid system. We emphasize that we observe a finite squeezing level in a comparatively broad pump-power range and thus even in the case of no (very weak) nondegenerate gain, as also shown in other experimental works, such as Ref. [61].

APPENDIX E: QUADRATURE VARIANCES OF THE SPIN MODE

To determine the quadrature variances of the spin mode and calculate the corresponding squeezing level when probing the electron-spin ensemble with squeezed microwave signals, we rely on the steady-state solution of the input-output equations. As presented in Eq. (4), the signal variances of the experimentally accessible input and output microwave fields are not directly dependent on the variances of the spin ensemble, and consequently, we have to extend our theoretical model in Sec. III.

In the following model, we rely on the theoretical description of the spin ensemble as a harmonic oscillator [22,34,35]. The spin ensemble is described via the bosonic annihilation and creation operators, $\hat{s}$ and $\hat{s}^\dagger$, respectively. The corresponding spin mode is characterized by a finite linewidth γ_s and resonance frequency ω_s . Within this model, we can define the quadratures of the spin excitation as

$$\hat{q}_s(t) = \frac{1}{2}[\hat{s}(t) + \hat{s}^\dagger(t)], \quad \hat{p}_s(t) = \frac{1}{2i}[\hat{s}(t) - \hat{s}^\dagger(t)], \quad (\text{E1})$$

which formally correspond to definitions in the frequency domain that read

$$\begin{aligned} \hat{q}_s(\omega) &= \frac{1}{\sqrt{2}}[\hat{s}(\omega) + \hat{s}^\dagger(-\omega)], \\ \hat{p}_s(\omega) &= \frac{1}{\sqrt{2}i}[\hat{s}(\omega) - \hat{s}^\dagger(-\omega)]. \end{aligned} \quad (\text{E2})$$

Following the ideas in [45,62], we can derive an effective rate model for the quadrature variances. We have to assume a regime where the resonator dynamics is fast compared to the intrinsic spin dephasing in order to adiabatically eliminate the cavity field and treat it as an effective reservoir for the collective spin excitation [46,47]. Within this model, we consider that the spin mode interacts with two independent reservoirs: (i) a cavity-induced squeezed bath characterized by effective noise parameters N_{eff} and M_{eff} and (ii) its intrinsic thermal bath with rate γ_s . The quadrature variances of the harmonic-oscillator approximation of the spin ensemble then follow as

$$\begin{aligned} \sigma_{\hat{q}_s}^2 &= \frac{1}{4}(1 + 2N_{s,\text{eff}} - 2|M_{s,\text{eff}}|), \\ \sigma_{\hat{p}_s}^2 &= \frac{1}{4}(1 + 2N_{s,\text{eff}} + 2|M_{s,\text{eff}}|). \end{aligned} \quad (\text{E3})$$

The corresponding spin noise parameters are defined as

$$N_{s,\text{eff}} = \frac{\Gamma_c N_{\text{eff}} + \gamma_s \bar{n}_s^{\text{th}}}{\Gamma_c + \gamma_s}, \quad M_{s,\text{eff}} = \frac{\Gamma_c M_{\text{eff}}}{\Gamma_c + \gamma_s}, \quad (\text{E4})$$

where $\bar{n}_s^{\text{th}}$ is defined as the photon number of the internal spin bath causing spin dephasing, $r = S/[20 \log_{10}(e)]$ is the squeezing factor, and $\Gamma_c = 2g_{\text{eff}}^2/\kappa$ can be considered a Purcell-like damping rate. The noise parameters of the cavity-induced squeezed bath read

$$\begin{aligned} N_{\text{eff}} &= \frac{\kappa_{\text{ext}} \sinh(r)^2 + \kappa_{\text{int}} \bar{n}_{\text{th}}}{\kappa}, \\ M_{\text{eff}} &= \frac{\kappa_{\text{ext}} \sinh(r)^2 \cosh(r)}{\kappa}. \end{aligned} \quad (\text{E5})$$

This effective rate description provides a way to analytically determine the squeezing level of the spin mode depending on

the squeezing level of the input-field mode while considering direct and indirect couplings to an internal spin and resonator bath, respectively. We want to emphasize that this simplified model does not capture any time-dependent dynamics and

reflects only the steady state. Describing the transient regime of the considered hybrid system requires a more complex formalism of Lindblad master equations, which goes beyond the scope of the current work.

-
- [1] N. Gisin and R. Thew, Quantum communication, *Nat. Photon.* **1**, 165 (2007).
- [2] H. J. Kimble, The quantum internet, *Nature (London)* **453**, 1023 (2008).
- [3] M. K. Bhaskar, R. Riedinger, B. Machielse, D. S. Levonian, C. T. Nguyen, E. N. Knall, H. Park, D. Englund, M. Lončar, D. D. Sukachev, and M. Lukin, Experimental demonstration of memory-enhanced quantum communication, *Nature (London)* **580**, 60 (2020).
- [4] M. Teller, S. Plascencia, C. S. Jachimska, S. Grandi, and H. de Riedmatten, A solid-state temporally multiplexed quantum memory array at the single-photon level, *npj Quantum Inf.* **11**, 92 (2025).
- [5] W. Zhang, T. van Leent, K. Redeker, R. Garthoff, R. Schwonnek, F. Fertig, S. Eppelt, W. Rosenfeld, V. Scarani, C. C.-W. Lim, and H. Weinfurter, A device-independent quantum key distribution system for distant users, *Nature (London)* **607**, 687 (2022).
- [6] D. P. Nadlinger, P. Drmota, B. C. Nichol, G. Araneda, D. Main, R. Srinivas, D. M. Lucas, C. J. Ballance, K. Ivanov, E.-Z. Tan, R. Sekatski, R. L. Urbanke, R. Renner, N. Sangouard, and J.-D. Bancal, Experimental quantum key distribution certified by Bell's theorem, *Nature (London)* **607**, 682 (2022).
- [7] Y.-F. Pu, S. Zhang, Y.-K. Wu, N. Jiang, W. Chang, C. Li, and L.-M. Duan, Experimental demonstration of memory-enhanced scaling for entanglement connection of quantum repeater segments, *Nat. Photon.* **15**, 374 (2021).
- [8] S. Wiesner, Conjugate coding, *SIGACT News* **15**, 78 (1983).
- [9] F. Pastawski, N. Y. Yao, L. Jiang, M. D. Lukin, and J. I. Cirac, Unforgeable noise-tolerant quantum tokens, *Proc. Natl. Acad. Sci. USA* **109**, 16079 (2012).
- [10] K. Jiráková, K. Bartkiewicz, A. Černoč, and K. Lemr, Experimentally attacking quantum money schemes based on quantum retrieval games, *Sci. Rep.* **9**, 16318 (2019).
- [11] H. P. Specht, C. Nölleke, A. Reiserer, M. Uphoff, E. Figueroa, S. Ritter, and G. Rempe, A single-atom quantum memory, *Nature (London)* **473**, 190 (2011).
- [12] J.-S. Tang, Z.-Q. Zhou, Y.-T. Wang, Y.-L. Li, X. Liu, Y.-L. Hua, Y. Zou, S. Wang, D.-Y. He, G. Chen, Y.-N. Sun, Y. Yu, M.-F. Li, G.-W. Zha, H.-Q. Ni, Z.-C. Niu, C.-F. Li, and G.-C. Guo, Storage of multiple single-photon pulses emitted from a quantum dot in a solid-state quantum memory, *Nat. Commun.* **6**, 8652 (2015).
- [13] A. Wallucks, I. Marinković, B. Hensen, R. Stockill, and S. Gröblacher, A quantum memory at telecom wavelengths, *Nat. Phys.* **16**, 772 (2020).
- [14] N. Bar-Gill, L. M. Pham, A. Jarmola, D. Budker, and R. L. Walsworth, Solid-state electronic spin coherence time approaching one second, *Nat. Commun.* **4**, 1743 (2013).
- [15] M. Zhong, M. P. Hedges, R. L. Ahlefeldt, J. G. Bartholomew, S. E. Beavan, S. M. Wittig, J. J. Longdell, and M. J. Sellars, Optically addressable nuclear spins in a solid with a six-hour coherence time, *Nature (London)* **517**, 177 (2015).
- [16] A. Ghirri, C. Bonizzoni, F. Troiani, N. Buccheri, L. Beverina, A. Cassinese, and M. Affronte, Coherently coupling distinct spin ensembles through a high- T_c superconducting resonator, *Phys. Rev. A* **93**, 063855 (2016).
- [17] A. Strinić, P. Oehrl, A. Marx, P. A. Bushev, H. Huebl, R. Gross, and N. Kukharchyk, Broadband electron paramagnetic resonance spectroscopy of $^{167}\text{Er}:\text{LiYF}_4$ at millikelvin temperatures, *Phys. Rev. B* **111**, 214430 (2025).
- [18] N. Timoney, I. Usmani, P. Jobez, M. Afzelius, and N. Gisin, Single-photon-level optical storage in a solid-state spin-wave memory, *Phys. Rev. A* **88**, 022324 (2013).
- [19] C. Grezes, B. Julsgaard, Y. Kubo, W. L. Ma, M. Stern, A. Bienfait, K. Nakamura, J. Isoya, S. Onoda, T. Ohshima, V. Jacques, D. Vion, D. Esteve, R. B. Liu, K. Mølmer, and P. Bertet, Storage and retrieval of microwave fields at the single-photon level in a spin ensemble, *Phys. Rev. A* **92**, 020301(R) (2015).
- [20] J. O'Sullivan, O. W. Kennedy, K. Debnath, J. Alexander, C. W. Zollitsch, M. Šimėnas, A. Hashim, C. N. Thomas, S. Withington, I. Siddiqi, K. Mølmer, and J. J. L. Morton, Random-access quantum memory using chirped pulse phase encoding, *Phys. Rev. X* **12**, 041014 (2022).
- [21] C. Eichler, A. J. Sigillito, S. A. Lyon, and J. R. Petta, Electron spin resonance at the level of 10^4 spins using low impedance superconducting resonators, *Phys. Rev. Lett.* **118**, 037701 (2017).
- [22] A. Bienfait, P. Campagne-Ibarcq, A. H. Kiilerich, X. Zhou, S. Probst, J. J. Pla, T. Schenkel, D. Vion, D. Esteve, J. J. L. Morton, K. Mølmer, and P. Bertet, Magnetic resonance with squeezed microwaves, *Phys. Rev. X* **7**, 041011 (2017).
- [23] L. Zhong, E. P. Menzel, R. Di Candia, P. Eder, M. Ihmig, A. Baust, M. Haeberlein, E. Hoffmann, K. Inomata, T. Yamamoto, Y. Nakamura, E. Solano, F. Deppe, A. Marx, and R. Gross, Squeezing with a flux-driven Josephson parametric amplifier, *New J. Phys.* **15**, 125013 (2013).
- [24] S. Pogorzalek, K. G. Fedorov, L. Zhong, J. Goetz, F. Wulschner, M. Fischer, P. Eder, E. Xie, K. Inomata, T. Yamamoto, Y. Nakamura, A. Marx, F. Deppe, and R. Gross, Hysteretic flux response and nondegenerate gain of flux-driven Josephson parametric amplifiers, *Phys. Rev. Appl.* **8**, 024012 (2017).
- [25] M. Renger, S. Pogorzalek, Q. Chen, Y. Nojiri, K. Inomata, Y. Nakamura, M. Partanen, A. Marx, R. Gross, F. Deppe, and K. G. Fedorov, Beyond the standard quantum limit for parametric amplification of broadband signals, *npj Quantum Inf.* **7**, 160 (2021).
- [26] K. E. Honasoge, M. Handschuh, W. K. Yam, S. Gandorfer, D. Bazulin, N. Bruckmoser, L. Koch, A. Marx, R. Gross, and K. G. Fedorov, Fabrication of low-loss Josephson parametric devices, *Phys. Rev. B* **111**, 214508 (2025).
- [27] E. P. Menzel, F. Deppe, M. Mariani, M. A. Araque Caballero, A. Baust, T. Niemczyk, E. Hoffmann, A. Marx,

- E. Solano, and R. Gross, Dual-path state reconstruction scheme for propagating quantum microwaves and detector noise tomography, *Phys. Rev. Lett.* **105**, 100401 (2010).
- [28] M. Mariantoni, E. P. Menzel, F. Deppe, M. A. Araque Caballero, A. Baust, T. Niemczyk, E. Hoffmann, E. Solano, A. Marx, and R. Gross, Planck spectroscopy and quantum noise of microwave beam splitters, *Phys. Rev. Lett.* **105**, 133601 (2010).
- [29] S. Gandorfer, M. Renger, W. K. Yam, F. Fesquet, A. Marx, R. Gross, and K. G. Fedorov, Two-dimensional Planck spectroscopy for microwave photon calibration, *Phys. Rev. Appl.* **23**, 024064 (2025).
- [30] K. G. Fedorov, S. Pogorzalek, U. Las Heras, M. Sanz, P. Yard, P. Eder, M. Fischer, J. Goetz, E. Xie, K. Inomata, Y. Nakamura, R. Di Candia, E. Solano, A. Marx, F. Deppe, and R. Gross, Finite-time quantum entanglement in propagating squeezed microwaves, *Sci. Rep.* **8**, 6416 (2018).
- [31] S. Weichselbaumer, M. Zens, C. W. Zollitsch, M. S. Brandt, S. Rotter, R. Gross, and H. Huebl, Echo trains in pulsed electron spin resonance of a strongly coupled spin ensemble, *Phys. Rev. Lett.* **125**, 137701 (2020).
- [32] J. Gordon and K. Bowers, Microwave spin echoes from donor electrons in silicon, *Phys. Rev. Lett.* **1**, 368 (1958).
- [33] G. Feher, Electron spin resonance experiments on donors in silicon. I. Electronic structure of donors by the electron nuclear double resonance technique, *Phys. Rev.* **114**, 1219 (1959).
- [34] J. H. Wesenberg, A. Ardavan, G. A. D. Briggs, J. J. L. Morton, R. J. Schoelkopf, D. I. Schuster, and K. Mølmer, Quantum computing with an electron spin ensemble, *Phys. Rev. Lett.* **103**, 070502 (2009).
- [35] I. Chiorescu, N. Groll, S. Bertaina, T. Mori, and S. Miyashita, Magnetic strong coupling in a spin-photon system and transition to classical regime, *Phys. Rev. B* **82**, 024413 (2010).
- [36] I. Diniz, S. Portolan, R. Ferreira, J. M. Gérard, P. Bertet, and A. Auffeves, Strongly coupling a cavity to inhomogeneous ensembles of emitters: Potential for long-lived solid-state quantum memories, *Phys. Rev. A* **84**, 063810 (2011).
- [37] F. Fesquet, Demonstration of microwave single-shot quantum key distribution, Ph.D. thesis, Technische Universität München, 2025, https://www.wmi.badw.de/fileadmin/WMI/Publications/Fesquet_Florian_Doktorarbeit_2025.pdf.
- [38] C. W. Gardiner and M. J. Collett, Input and output in damped quantum systems: Quantum stochastic differential equations and the master equation, *Phys. Rev. A* **31**, 3761 (1985).
- [39] D. A. Steck, *Quantum and Atom Optics* (2023).
- [40] K. Sandner, H. Ritsch, R. Amsüss, C. Koller, T. Nöbauer, S. Putz, J. Schmiedmayer, and J. Majer, Strong magnetic coupling of an inhomogeneous nitrogen-vacancy ensemble to a cavity, *Phys. Rev. A* **85**, 053806 (2012).
- [41] M. Aspelmeyer, T. J. Kippenberg, and F. Marquardt, Cavity optomechanics, *Rev. Mod. Phys.* **86**, 1391 (2014).
- [42] S. Probst, F. Song, P. A. Bushev, A. V. Ustinov, and M. Weides, Efficient and robust analysis of complex scattering data under noise in microwave resonators, *Rev. Sci. Instrum.* **86**, 024706 (2015).
- [43] S. L. Braunstein and P. van Loock, Quantum information with continuous variables, *Rev. Mod. Phys.* **77**, 513 (2005).
- [44] M. Müller, T. Luschmann, A. Faltermeier, S. Weichselbaumer, L. Koch, G. B. P. Huber, H. W. Schumacher, N. Ubbelohde, D. Reifert, T. Scheller, F. Deppe, A. Marx, S. Filipp, M. Althammer, R. Gross, and H. Huebl, Magnetic field robust high quality factor NbTiN superconducting microwave resonators, *Mater. Quantum Technol.* **2**, 015002 (2022).
- [45] D. F. Walls and G. J. Milburn, *Quantum Optics*, 1st ed. (Springer, Berlin, Heidelberg, 1994).
- [46] F. Damanet, A. J. Daley, and J. Keeling, Atom-only descriptions of the driven-dissipative Dicke model, *Phys. Rev. A* **99**, 033845 (2019).
- [47] S. Giaccari, G. Dellea, M. G. Genoni, and G. Bertaina, Higher-order adiabatic elimination in atom-cavity systems and its impact on spin-squeezing generation, *Quantum Sci. Technol.* **11**, 035019 (2026).
- [48] M. Afzelius, N. Sangouard, G. Johansson, M. Staudt, and C. Wilson, Proposal for a coherent quantum memory for propagating microwave photons, *New J. Phys.* **15**, 065008 (2013).
- [49] B. Julsgaard, C. Grezes, P. Bertet, and K. Mølmer, Quantum memory for microwave photons in an inhomogeneously broadened spin ensemble, *Phys. Rev. Lett.* **110**, 250503 (2013).
- [50] L. Greggio, T. Lorriaux, A. Petrescu, M. Mirrahimi, and A. Bienfait, Optimal absorption and emission of itinerant fields into a spin-ensemble memory, *Phys. Rev.* (2026), doi:10.1103/6xjr-tflh.
- [51] J. Z. Bernád, M. Schilling, Y. Wen, M. M. Müller, T. Calarco, P. Bertet, and F. Motzoi, Analytical solutions for optimal photon absorption into inhomogeneous spin memories, *J. Phys. B* **58**, 035501 (2025).
- [52] F. Fesquet, F. Kronowetter, M. Renger, W. K. Yam, S. Gandorfer, K. Inomata, Y. Nakamura, A. Marx, R. Gross, and K. G. Fedorov, Demonstration of microwave single-shot quantum key distribution, *Nat. Commun.* **15**, 7544 (2024).
- [53] P. Oehrl and H. Huebl, Probing an electron spin ensemble with squeezed microwave signals, [Dataset], Zenodo, 2026, <https://doi.org/10.5281/zenodo.21260844>.
- [54] A. Schweiger and G. Jeschke, *Principles of Pulse Electron Paramagnetic Resonance* (Oxford University Press, Oxford, 2001).
- [55] S. Weichselbaumer, P. Natzkin, C. W. Zollitsch, M. Weiler, R. Gross, and H. Huebl, Quantitative modeling of superconducting planar resonators for electron spin resonance, *Phys. Rev. Appl.* **12**, 024021 (2019).
- [56] C. W. Zollitsch, K. Mueller, D. P. Franke, S. T. B. Goennenwein, M. S. Brandt, R. Gross, and H. Huebl, High cooperativity coupling between a phosphorus donor spin ensemble and a superconducting microwave resonator, *Appl. Phys. Lett.* **107**, 142105 (2015).
- [57] E. L. Hahn, Spin echoes, *Phys. Rev.* **80**, 580 (1950).
- [58] A. M. Tyryshkin, S. A. Lyon, A. V. Astashkin, and A. M. Raitsimring, Electron spin relaxation times of phosphorus donors in silicon, *Phys. Rev. B* **68**, 193207 (2003).
- [59] C. Eichler, D. Bozyigit, C. Lang, M. Baur, L. Steffen, J. M. Fink, S. Filipp, and A. Wallraff, Observation of two-mode

- squeezing in the microwave frequency domain, [Phys. Rev. Lett. **107**, 113601 \(2011\)](#).
- [60] S. Boutin, D. M. Toyli, A. V. Venkatramani, A. W. Eddins, I. Siddiqi, and A. Blais, Effect of higher-order nonlinearities on amplification and squeezing in Josephson parametric amplifiers, [Phys. Rev. Appl. **8**, 054030 \(2017\)](#).
- [61] K. G. Fedorov, L. Zhong, S. Pogorzalek, P. Eder, M. Fischer, J. Goetz, E. Xie, F. Wulschner, K. Inomata, T. Yamamoto, Y. Nakamura, R. Di Candia, U. L. Heras, M. Sanz, E. Solano, E. P. Menzel, F. Deppe, A. Marx, and R. Gross, Displacement of propagating squeezed microwave states, [Phys. Rev. Lett. **117**, 020502 \(2016\)](#).
- [62] J. Wang, H. M. Wiseman, and G. J. Milburn, Dynamical creation of entanglement by homodyne-mediated feedback, [Phys. Rev. A **71**, 042309 \(2005\)](#).