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Nb-Ti-N nanowire resonators and prospects for spin-photon coupling with electrons on solid neon

Y. Tian¹, I. Grytsenko¹, A. Jennings¹, J. Wang¹, H. Ikegami¹, X. Zhou^{1,2}, S. Tamate¹, H. Terai³, H. Kutsuma⁴, D. Jin⁵, M. Benito^{6,7} and E. Kawakami^{1,8,*}

¹*RIKEN Center for Quantum Computing, 2-1 Hirosawa, Wako, Saitama 351-0198, Japan*

²*Department of Mechanical Engineering, FAMU-FSU College of Engineering, Florida State University, Tallahassee, Florida 32310, USA*

³*Advanced ICT Research Institute, National Institute of Information and Communications Technology (NICT), Kobe 651-2492, Japan*

⁴*Graduate School of Engineering, Tohoku University, Sendai 980-8579, Japan*

⁵*Department of Physics and Astronomy, University of Notre Dame, Notre Dame, Indiana 46556, USA*

⁶*Institute of Physics, University of Augsburg, Augsburg 86159, Germany*

⁷*Center for Advanced Analytics and Predictive Sciences, University of Augsburg, 86135 Augsburg, Germany*

⁸*RIKEN Cluster for Pioneering Research, 2-1 Hirosawa, Wako, Saitama 351-0198, Japan*



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Electrons floating on solid neon exhibit long charge-coherence times, making them attractive for hybrid quantum systems. When combined with high-quality high-impedance superconducting resonators and a local magnetic field gradient, this platform enables strong charge-photon and spin-charge coupling—key ingredients for scalable spin-qubit architectures. In this work, we demonstrate that Nb-Ti-N nanowire resonators maintain high quality factors around 10^5 after depositing solid neon onto the resonators and subsequently loading electrons onto the neon surface, validating their suitability for electrons-on-neon platforms. Building on these experimental results, we theoretically analyze micromagnet designs and coupling strategies that can enable spin-photon interactions in this platform. Our analysis outlines performance targets for next-generation devices, showing that, at the charge sweet spot, spin-qubit gate fidelities exceeding 99.99% for single-qubit operations and 99.9% for two-qubit operations are achievable with natural neon.

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I. INTRODUCTION

Using cavity quantum electrodynamics (cQED), the quantum states of an electron can be coupled to a microwave photon in a superconducting resonator. This approach has been successfully implemented using electrons in semiconductors, for both charge [1–3] and spin [4–6] degrees of freedom. Recent experiments have demonstrated coupling between distant qubits [7–9] and, more recently, realized remote two-qubit gate operations [10], contributing to the development of scalable electron spin-qubit systems.

Recently, cQED experiments have been demonstrated using electrons on solid neon [11]. The development of an electron charge qubit on solid neon is a significant breakthrough, reaching charge coherence times of about 100 μ s [11,12]. Recent theoretical studies suggest that at temperatures around approximately 10 mK, decoherence of the

charge degree of freedom of the electron is primarily limited by acoustic phonons in solid neon [13]. However, the coherence time expected from this mechanism is longer than the measured value, which indicates that the observed decoherence may instead be dominated by stray electrons near the qubit [14]. Therefore, controlling the electron density on solid neon is a key step toward achieving qubits with longer coherence times.

Even longer coherence times, potentially reaching the order of 1 s, are expected for the spin states of an electron on solid neon [15], making them attractive for spin-photon coupling schemes. Such coupling, mediated by charge-photon interaction, has been proposed [16,17] and experimentally demonstrated with electrons in semiconductors [6,10,18]. In this framework, maintaining long coherence times for both spin and charge states is crucial for achieving high spin-qubit gate fidelities.

By combining cQED technology with these extended coherence times for both the charge and spin states of electrons on solid neon, electron qubit systems can be scaled

*Contact author: e2006k@gmail.com

into large networks, while simultaneously simplifying the complexity of control required for quantum error correction [19]. Along this line, the primary goal is to achieve strong coupling between the spin state of a single electron on solid neon and a microwave photon (spin-photon strong coupling). In order to realize spin-photon strong coupling, spin-charge interaction is required. The charge state can couple to a photon in a superconducting resonator, enabling spin-photon coupling via the spin-charge interaction. This spin-charge interaction can be artificially introduced using locally placed microferromagnets [20–23]. An external magnetic field is applied to magnetize the ferromagnet and induce a finite Zeeman splitting.

As a first step toward realizing spin qubits using electrons on neon and a nanowire resonator, we fabricated Nb-Ti-N superconducting nanowire resonators without integrating micromagnets and carried out foundational experiments. In these experiments, we deposited a thin layer of neon followed by electrons onto the resonator, gradually increased the electron density, and monitored the resulting resonance peak shift. Importantly, depositing neon and electrons did not degrade the internal quality factor of the resonators. Encouraged by these experimental results, we proceeded to a theoretical analysis of the next stage toward implementing spin qubits, using the experimentally obtained values as input. We investigated the optimized geometry for micromagnets required to induce local magnetic field gradients, and demonstrated that cooperativity between the resonator photons and the spin state can reach $\gtrsim 10^7$. We theoretically identified optimal conditions that enable high-fidelity single- and two-qubit gate operations for spin qubits, and found that fidelities of 99.99% and 99.9%, respectively, are potentially achievable even when using natural neon.

II. NANOWIRE RESONATOR

Since applying a magnetic field is necessary, we fabricated resonators using a 20-nm-thick Nb-Ti-N film sputtered onto high-resistivity silicon to enhance their resilience to magnetic fields. In addition, Nb-Ti-N has high kinetic inductance, which, in turn, strengthens the coupling between the photons of the resonator and the charge state [5]. The critical temperature of the Nb-Ti-N film was measured to be 10.7 K. Achieving strong coupling between the photons of the resonator and the charge state is crucial for obtaining high spin-photon cooperativity, which is essential for realizing high spin-qubit gate fidelities [10,21]. To further improve performance, a nanowire-shaped resonator was chosen to suppress the generation of vortices that could disrupt superconductivity [5,24].

We fabricated three resonators, all sharing a single feed-line (Fig. 1). The two ends of each resonator meet with a small gap, where the electric field is maximized, enabling

strong coupling between the dipole moment of the electron and the electric field. Both the widths and the gaps were designed to be 100 nm for resonator 1, 200 nm for resonator 2, and 300 nm for resonator 3. The resonance frequency of resonator 1 was measured to be $f_r = 4.81$ GHz with $Q_{\text{ext}} = 3.7 \times 10^4$, showing good agreement with simulations (see Appendix B 2 and Table I). $Q_{\text{int}} = 2.3 \times 10^5$ and $Q_{\text{tot}} = 3.2 \times 10^4$ were measured at 10 mK. These values were obtained using an input power of -21 dBm, corresponding to an estimated intracavity average photon number of $\langle n_{\text{ph}} \rangle \approx 10^5$. The resonance frequency of resonator 2 was measured to be $f_r = 5.91$ GHz. Although both the width and the gap were nominally designed to be 200 nm, over-etching appears to have reduced the resonator width, resulting in better agreement with simulation results for a width-gap configuration in the range of 150–150 nm (see Table I and Appendix B 2). Other parameters for resonators 2 and 3 are provided in Table I. For all the resonators, the total loss rate $\kappa/2\pi = f_r/Q_{\text{tot}} \approx 0.1$ MHz.

We note that this device design intentionally omitted dc bias lines to avoid degrading the Q_{int} of the Nb-Ti-N resonator, as the primary objective of this work was to establish a performance baseline of the resonators with solid neon and electrons.

III. NEON AND ELECTRON DEPOSITION

We deposited neon at 25 K. Based on the resonance-frequency shifts of resonators 1 and 2, measured at 10 mK as -0.94% and -0.86% , respectively, we estimated that the neon thicknesses in the regions of resonators 1 and 2 were 160 nm and 270 nm at most, respectively (see Appendix B). From the volume of neon gas recovered after warming up, we estimated that the total amount of neon deposited in the experimental cell was 0.2 mol. If we assume that the neon was deposited uniformly from the bottom of the experimental cell (surface area = 14 cm^2), the expected neon thickness would be 1 mm. This estimate is inconsistent with the resonance-frequency shifts, suggesting that neon is primarily deposited on the walls of the experimental cell and capillary lines. This spread of neon to these surfaces may occur due to triple-point wetting during the passage through the triple point when depositing solid neon [27,28]. Aside from the thickness estimation, we also found that compared to the bare resonator, Q_{int} measured at 10 mK and at low photon number is increased with the presence of neon (see Appendix C). This must be attributed to the fact that the higher dielectric constant of neon reduces the participation ratio in the region in which the two-level-system (TLS) exists.

After depositing neon, we proceeded with electron deposition by sending a voltage pulse to a filament placed in the vacuum cell. In Fig. 2(a), we show the resonance-frequency shifts measured at 3.4 K by increasing the height

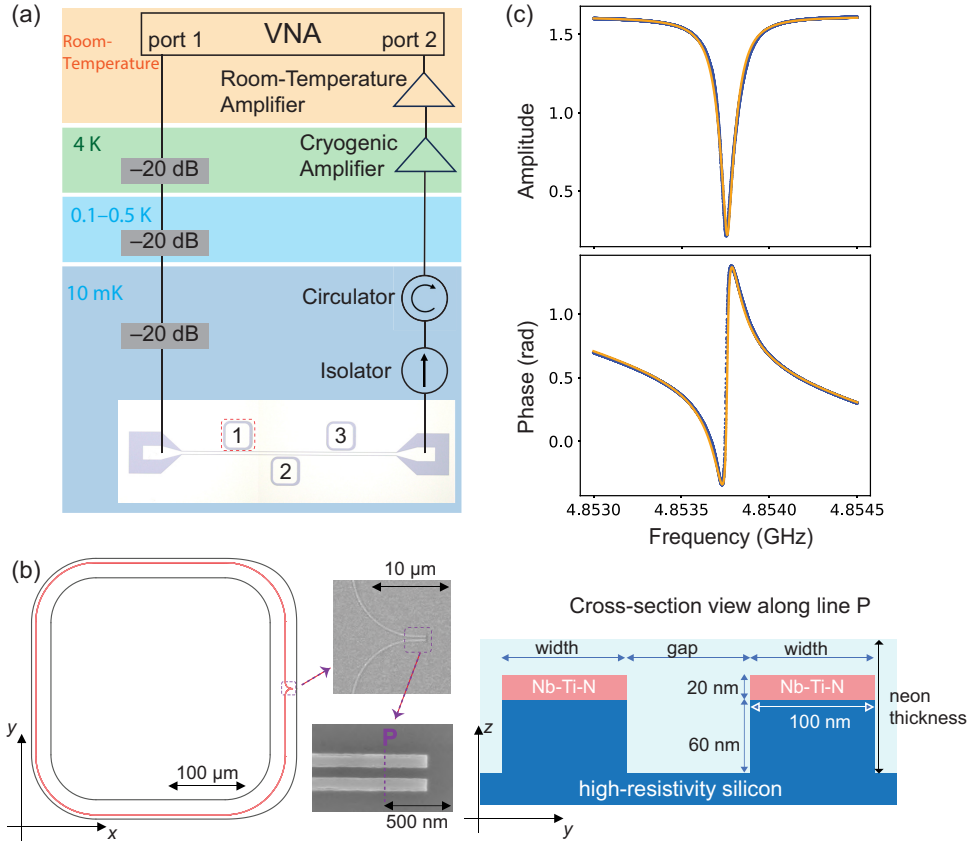


FIG. 1. (a) A schematic overview of the electrical setup within the dilution refrigerator and at room temperature. The input lines are attenuated with attenuators by 60 dB in total, with the output lines equipped with an LNF-LNC4_8C amplifier at 4 K and an LNF-LNR4_8ART amplifier at room temperature. The diagram also includes an optical micrograph of the measured device on the 10-mK plate, showing three resonators—resonators 1, 2, and 3—all sharing a single microwave (MW) feedline. (b) A conceptual illustration of resonator 1. The nanowire resonator is shown in pink. The two rounded black squares represent grounded regions, with both the interior of the squares and the surrounding exterior connected to ground. Scanning-electron-microscopy images of the enlarged areas highlight the two ends of the resonator. A schematic illustration of vertical cross sections along line P is also presented. (c) The transmission response of resonator 1 measured at 10 mK with a vector network analyzer (VNA). The blue lines represent the measured S_{21} in linear amplitude and phase before the deposition of solid neon or electrons. The orange lines are constructed from the best fit using the method described in Refs. [25,26].

of the voltage pulse (duration approximately 3 s) applied to the filament. As shown in Fig. 2(b), the resonance frequency decreases with increasing electron number, while no change is observed in Q_{int} , as shown in Fig. 2(c). The maximum measured frequency shift was -0.9% .

The decrease in resonance frequency with increasing electron number on the neon surface is attributed to the surface conductivity of the two-dimensional (2D) electron layer. To illustrate this, we first consider the Drude model: $\sigma^{2D} = (e^2 n_e \tau / m_e) 1 / (1 + i\omega\tau)$, where e is the elementary charge, n_e is the surface electron density, τ is the scattering time, m_e is the electron mass, ω is the angular frequency of the microwave signal, and i is the imaginary unit. The highest reported τ for electrons on solid neon is 4.7 ps, based on low-frequency mobility measurements at 4.2 K [29,30]. A shorter value of 1.9 ps has also been reported for rougher surfaces [29].

We simulated the resonance-frequency shift Δf and the electron-induced loss $1/Q_e$, incorporating surface conductivity based on the Drude model (see Appendix B 2 b). The simulations show that a shift of -0.9% corresponds to $n_e = 0.4(0.8) \times 10^9 \text{ cm}^{-2}$ and $1/Q_e = 4.0(3.9) \times 10^{-3}$ for $\tau = 4.7(1.9)$ ps. Since τ is not directly measured and the Drude model does not accurately capture the measured Q_{int} , as discussed below, these values of n_e should be interpreted as indicative. The total internal quality factor is given by $1/Q_{\text{int}} = 1/Q_{\text{int},0} + 1/Q_e$, where $Q_{\text{int},0}$ is the value without electrons. The simulated $1/Q_{\text{int}}$ is nearly an order of magnitude larger than the experimentally measured value, which remains around 4×10^{-4} at 3.4 K and shows no dependence on electron density [see Fig. 2(c)]. This discrepancy suggests that the Drude model fails to capture the measured Q_{int} . In Appendix B 2 c, we explore a model that incorporates surface disorder via harmonic

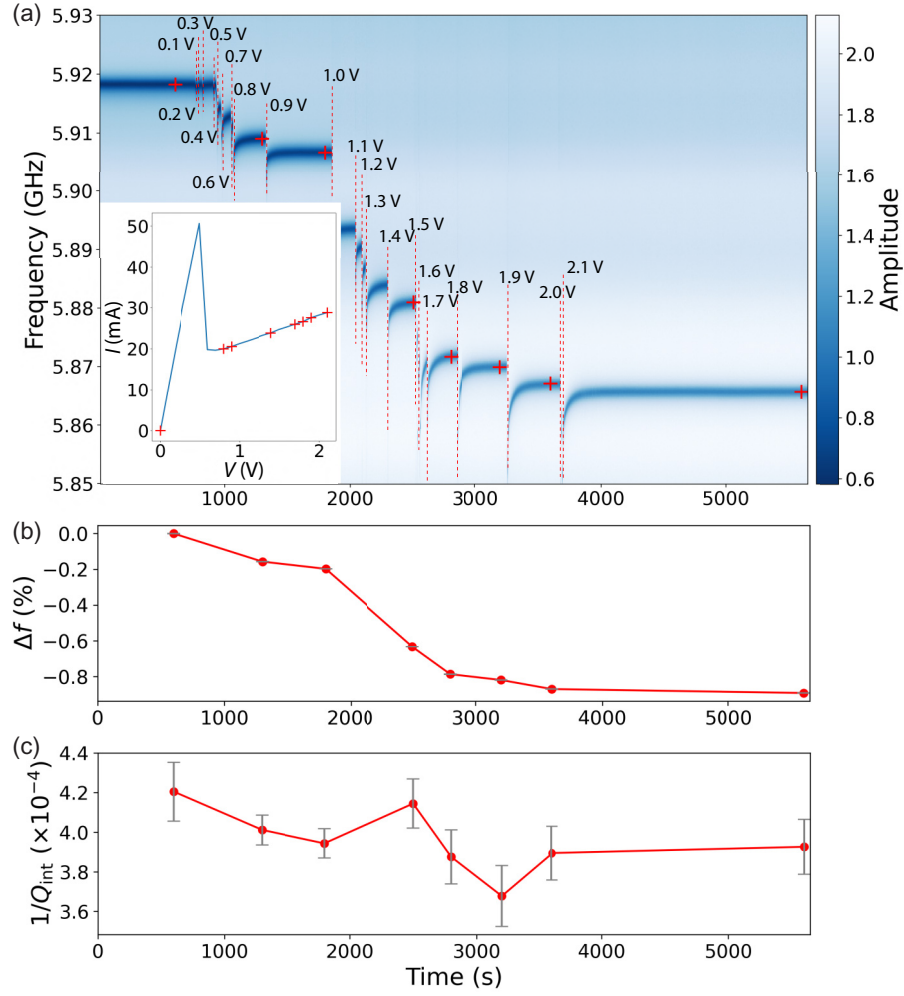


FIG. 2. (a) The resonance peaks of resonator 2 measured at 3.4 K over time. The vertical dashed red lines indicate the moments when electrons were emitted, with the applied filament voltages shown next to each. The red crosses mark the resonance peaks used to extract the resonance-frequency shift and Q_{int} in (b) and (c), respectively. The inset displays the I - V curve of the filament. (b) The resonance-frequency shift Δf , where the errors are smaller than the data points. (c) The inverse quality factor $1/Q_{\text{int}}$. The error bars correspond to 95% confidence.

TABLE I. The COMSOL simulations and experimental data measured at 10 mK without neon or electrons. The intended width-gap dimensions for resonators 1, 2, and 3 were 100 nm-100 nm, 200 nm-200 nm, and 300 nm-300 nm, respectively. For resonator 2, simulation results for width-gap dimensions of 150 nm-150 nm, which are closer to those of the actual fabricated sample, are also included. For more details of the simulations, see Appendix B 1. For the experimental data, a power of -21 dBm generated from the vector network analyzer (VNA) was used, corresponding to an estimated photon number in the resonator of $\langle n \rangle_{\text{ph}} \approx 10^5$. For power-dependence measurements, see Appendix C.

Resonator	Width-gap (nm)	Length (μm)	COMSOL simulation		Experimental data			
			f_r (GHz)	Q_{ext}	f_r (GHz)	Q_{int}	Q_{ext}	Q_{tot}
1	100-100	1450	4.72	4.0×10^4	4.81	2.3×10^5	3.7×10^4	3.2×10^4
2	200-200	1450	6.42	2.7×10^4	5.91	1.8×10^5	2.5×10^4	2.2×10^4
	150-150	1450	5.66	3.3×10^4				
3	300-300	1450	7.66	2.0×10^4	6.98	1.7×10^5	3.8×10^5	1.2×10^5

confinement, which can localize electrons. The results show that in certain parameter regimes, this approach can account for the observed frequency shift without a significant reduction in Q_e . These findings suggest that the limitations of the Drude model may arise from localization effects induced by surface roughness on a length scale smaller than the scattering length. In addition, the difference in the neon thickness observed between the two resonators can be attributed to large-scale nonuniformity of the neon deposition across the chip.

After depositing electrons at 3.4 K, the sample was cooled down to 10 mK, where two-tone spectroscopy was performed (see Appendix D) to confirm that the electrons remained trapped during the cooling process.

IV. MAGNETIC FIELD GRADIENT

The sample used in this work consists exclusively of resonators. As depicted in Fig. 3, we propose to add two Co/Ti/(Nb-Ti-N) electrodes at both ends of resonator 1, on its upper and lower sides. Cobalt (Co) serves as a ferromagnetic layer, titanium (Ti) as an adhesion layer, and Nb-Ti-N as part of the resonator. These electrodes are designed

to serve three roles: generating local magnetic field gradients, tuning the confinement potential of the electron through the application of dc voltage to the electrodes, and performing electric dipole spin resonance (EDSR) [10,20,31,32] by applying microwave (MW) signals to one of the electrodes (Sec. VI).

In the charge-qubit experiments reported in Refs. [11,12], a single electron trapped on solid neon shows a spectrum that qualitatively equals that seen in semiconductor double quantum dots (i.e., the energy of the charge state follows $\sqrt{\epsilon^2 + 4t_c^2}$, where ϵ is the charge detuning and t_c is the effective tunnel coupling [21]), which may be due to the surface roughness of the solid neon [33]. Here, we assume that the two minima of the electrical potential are created at the edges of each end of the resonator. This assumption is plausible because these points are where the attraction toward the surface becomes maximum due to the image potential created by the resonators. This configuration is illustrated in Fig. 3, and the distance between the two minima, i.e., the interdot distance $d = 100$ nm.

By applying an external magnetic field B_{ext} along the y axis, the Co parts are magnetized along the same axis and induce local magnetic fields. We present the

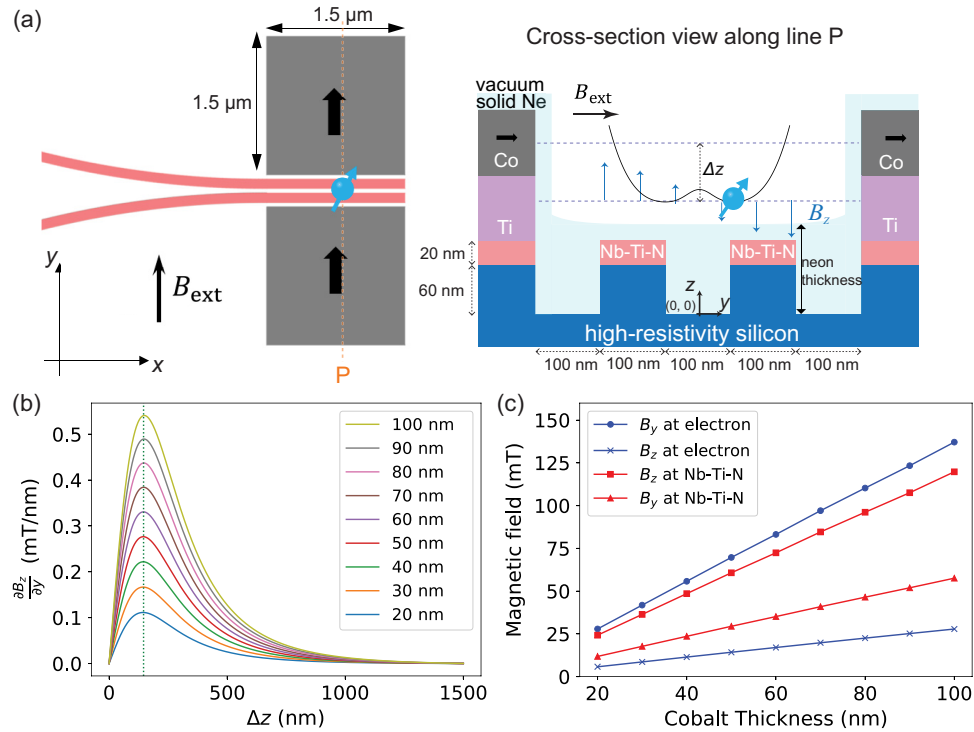


FIG. 3. (a) The addition of Co magnets (gray) near the gap between the two ends of resonator 1. In the cross-section view, the black line schematically represents the electrical potential experienced by the electron. (b) The magnetic field gradient in the z direction along the y axis, $\partial B_z / \partial y$, is shown as a function of the vertical distance Δz between the center position of the Co magnet and the electron, for various Co thicknesses ranging from 20 nm to 100 nm. (c) The absolute values of the offsets of the y and z components of the magnetic field at the position of the electron (i.e., at the potential minima), as well as the y and z components at the position of the Nb-Ti-N resonator, are plotted as functions of the Co thickness. For each Co thickness, the Ti thickness is optimized to maximize $\partial B_z / \partial y$.

numerical-calculation results of the induced local magnetic fields and their gradients [34] in Fig. 3. Simulation details are given in Appendix E.

In Fig. 3(b), we show the z -direction magnetic field gradient along the y axis, $\partial B_z/\partial y$, as a function of the distance between the z -center position of the Co magnet and the position of the electron, Δz . We found that the z -direction magnetic field gradient becomes maximum when $\Delta z \approx 146$ nm [vertical dashed line in Fig. 3(b)]. For example, with a Co thickness of 65 nm, a Ti thickness of 196 nm (126 nm), a neon thickness of 160 nm (90 nm), and an electron floating height of 2.5 nm, $\Delta z = 146$ nm can be achieved. The value of Δz can be adjusted even after sample fabrication by modifying the neon thickness. With this configuration, we obtain $\partial B_z/\partial y = 0.36$ mT/nm, resulting in a magnetic field gradient in the z -direction between the two potential minima, expressed in radian frequency as $b_\perp = (g\mu_B/\hbar)(\partial B_z/\partial y)d = 2\pi \times 1$ GHz, where g is the free-electron g factor, μ_B is the Bohr magneton, and $\hbar$ is the reduced Planck constant. This magnetic field gradient strength is sufficient to reach the strong-coupling regime for spin-photon coupling, as discussed in Sec. V, and to achieve high-fidelity qubit gates, as discussed in Sec. VI.

In Fig. 3(c), we show the magnetic field offsets along the z and y axes introduced by the Co magnets at the position of the Nb-Ti-N resonators, which are at most $B_z = 60$ mT and $B_y = 130$ mT, respectively. These values are low enough not to degrade the quality factor of the resonators for a Nb-Ti-N thickness of 20 nm [5,24]. Small fabrication-induced displacements have little effect on this result (see Appendix E). The B_y experienced by the electron at the positions of the potential minima is $B_y = 90$ mT for the same configuration as discussed before (Co thickness of 65 nm and $\Delta z \approx 146$ nm). Thus, the total magnetic field that defines the Zeeman splitting in radian frequency is expressed as $b_\parallel \approx (g\mu_B/\hbar)(B_y + B_{\text{ext}})$. Thus, to set the Zeeman splitting equal to the resonator resonance frequency of 4.8 GHz, we need to apply $B_{\text{ext}} = 80$ mT, which is sufficient to magnetize Co and perform qubit measurements [9,10,18]. In the presence of a magnetic field gradient, charge noise—which causes fluctuations in the charge detuning ϵ —could lead to fluctuations in the spin splitting. This makes the spin qubit sensitive to charge noise. Here, two Co magnets are placed symmetrically with a narrow gap relative to the electron positions. As a result, the Zeeman splitting at the two potential minima becomes equal and the effect of charge noise at $\epsilon = 0$ is largely suppressed.

V. SPIN-PHOTON COUPLING AND LOSS RATE

In this section, we estimate the spin-photon coupling strength, evaluate the loss rates of the spin state and the resonator, and discuss the feasibility of reaching

the spin-photon strong-coupling regime. For this purpose, we use the parameters of resonator 1: an inductance of $L = 139$ nH, validated through experimental characterization (see Appendix F) and the measured resonance frequency $f_r = \omega_r/2\pi = 4.81$ GHz. The characteristic impedance of the resonator is then calculated as $Z_0 = 2f_r L = 1337$ Ω .

According to Ref. [5], the zero-point fluctuation (ZPF) of the rms voltage amplitude between the two ends of the resonator can be calculated as

$$V_0 = \frac{1}{\sqrt{2}} \frac{2L}{\pi} \sqrt{\frac{2\hbar\omega_r}{L}} \omega_r = 12 \mu\text{V} \quad (1)$$

and thus the charge-photon coupling strength can be calculated as

$$\frac{g_c}{2\pi} = \frac{\alpha e V_0 \cos \theta}{h}, \quad (2)$$

where $\theta = \arctan(\epsilon/2t_c)$ and $eV_0/h = 3$ GHz. Here, we use $\alpha = 0.05$. This yields $g_c/2\pi = 150$ MHz at $\epsilon = 0$, a value comparable to those measured in semiconductor quantum dots using Nb-Ti-N resonators with similar designs [10,18]. Moreover, COMSOL simulations indicate that α exceeds 0.1 for $d = 100$ nm, when the neon layer thickness is below 300 nm, suggesting that the chosen value does not lead to an overestimation.

The high-frequency noise-induced loss rate is given by $\gamma = \gamma_1/2 + \gamma_\phi$, where $\gamma_1 = 2\pi/T_1$, $\gamma_\phi = 2\pi/T_\phi$, T_1 is the energy relaxation time, and T_ϕ is the pure dephasing time. The quasistatic noise-induced loss rate is given in Ref. [35] $\gamma^* = 1/T_2^*$, where T_2^* is the inhomogeneous (quasistatic) dephasing time. We denote the loss rates for spin and charge by adding subscripts “s” and “c,” respectively, i.e., γ_s , $\gamma_{s,1}$, $\gamma_{s,\phi}$, and γ_s^* for spin, and γ_c , $\gamma_{c,1}$, $\gamma_{c,\phi}$, and γ_c^* for charge. The coherence times are denoted as $T_{c,1}$, $T_{c,\phi}$, $T_{c,2}^*$, and $T_{s,1}$, $T_{s,\phi}$, $T_{s,2}^*$, respectively.

To assess the impact of different loss mechanisms, we consider three representative scenarios, summarized in Table II and labeled as “thermal,” “^{nat}Ne,” and “²²Ne,” corresponding to scenarios 1, 2, and 3, respectively. Each scenario reflects a distinct set of assumptions regarding charge and spin loss rates. In *scenario 1* (“thermal”), we employed charge loss rates of intermediate magnitude, as experimentally reported in Ref. [13], representative of typical values found in previous studies. For the spin loss rate due to high-frequency noise, we used a value determined by thermal magnetic noise in copper. Although this effect is expected to be negligibly small—our electrodes are made of superconducting material, which exhibits minimal thermal magnetic noise, and both the printed circuit board and the experimental cell, though made of copper, are located far from the electrons—we include this case to evaluate a worst-case scenario. In *scenarios 2* (“^{nat}Ne”) and *scenario 3* (“²²Ne”), we used charge loss rates of intermediate magnitude, as experimentally reported in Ref. [13], representative of typical values found in previous studies. For the spin loss rate due to high-frequency noise, we used a value determined by thermal magnetic noise in copper. Although this effect is expected to be negligibly small—our electrodes are made of superconducting material, which exhibits minimal thermal magnetic noise, and both the printed circuit board and the experimental cell, though made of copper, are located far from the electrons—we include this case to evaluate a worst-case scenario.

and 3 (^{22}Ne), we employed the smallest experimentally measured charge loss rates reported in Ref. [12]. For the spin loss rates, we used values determined by hyperfine interactions with nuclear spins [15]. Scenario 2 corresponds to natural neon, whereas scenario 3 assumes isotopically purified neon containing only 1 ppm (part per million) of ^{21}Ne .

We set $\epsilon = 0$ from here onward for simplicity and to leverage the long coherence time at the charge sweet spot. To measure the strong coupling between the electron spin state and a microwave photon in the resonator, we tune the spin and photon energies into resonance by setting the Zeeman splitting equal to the resonator resonance frequency, i.e., the detuning of the spin state from the resonator $\Delta_s = b_{\parallel} - \omega_r = 0$. The spin-photon coupling is given by $g_s = \Lambda g_c$, where $\Lambda = \sin \tilde{\phi}$, in which $\tilde{\phi} = (\phi_+ + \phi_-)/2$ is the spin-charge mixing angle and $\phi_{\pm} = \arctan[b_{\perp}/(2t_c \pm b_{\parallel})]$ [36]. Due to spin-charge coupling, the effective high-frequency noise-induced and quasistatic spin loss rate at the charge sweet spot become [36]

$$\gamma'_s = \Lambda^2 \gamma_c + (1 - \Lambda^2) \gamma_s \quad (3)$$

and

$$\gamma_s^{*'} = \sqrt{\left(\Lambda^2 \frac{\gamma_c^*}{b_{\parallel}}\right)^2 + \left(\frac{(\cos \phi_+ + \cos \phi_-)}{2} \gamma_s^*\right)^2}, \quad (4)$$

respectively.

The condition for spectroscopically observing strong spin-photon coupling is $g_s \gg \gamma'_s, \gamma_s^{*'}, \kappa'$. Here, we examine whether this condition is satisfied for scenario 1. To evaluate the performance of the system under experimentally feasible conditions, we consider the following set of parameters as an example: $b_{\parallel}/2\pi = \omega_r/2\pi = 4.8$ GHz, $\epsilon = 0$, $2t_c/2\pi = 8.0$ GHz, and $b_{\perp}/2\pi = 1$ GHz. Under these conditions, we obtain $\Lambda = 0.19$, resulting in a spin-photon coupling strength of $g_s/2\pi = 28.5$ MHz. The total resonator decay rate, including Purcell decay induced by the coupling between the charge state and the resonator, is given by $\kappa' = \kappa + (g_c^2 \gamma_c / \Delta_c^2)$, where $\Delta_c = \sqrt{\epsilon^2 + 4t_c^2} - \omega_r$ is the detuning of the charge state from the resonator [16]. For the parameters considered here, $\kappa'/2\pi \approx \kappa/2\pi = 0.1$ MHz, as determined experimentally (see Sec. II). At the charge sweet spot, the effective quasistatic spin loss rate is not influenced by quasistatic charge noise. The spin loss rate due to high-frequency noise is calculated to be $\gamma'_s/2\pi \approx 7$ kHz. These parameters satisfy the condition for strong coupling. The corresponding cooperativity is estimated to be $C = g_s^2 / (\gamma'_s \kappa') \gtrsim 10^6$, confirming the realization of the strong spin-photon-coupling regime.

VI. QUBIT GATES

In this section, we investigate the optimal experimental conditions for both single-qubit gates and an entangling two-qubit gate. To realize single-qubit gates, we can employ EDSR [10,20,31,32] by applying MW signal to one of the Co/Ti/(Nb-Ti-N) electrodes and modulating the position of the electron along the y axis. During EDSR, we detune the resonator from both the spin and charge states to suppress decoherence through the resonator ($\Delta_c, \Delta_s \gg 0$). The EDSR Rabi frequency is determined by $f_r^s = \Lambda f_r^c$, where f_r^c is the charge Rabi frequency. Here, we use $f_r^c = 10$ MHz, as reported in Ref. [11]. In Ref. [11], the microwave for charge state operation is applied through the resonator, whereas in our case, it will be supplied via one of the Co/Ti/(Nb-Ti-N) electrodes, which are independent from the resonator. Although this difference could lead to a higher f_r^c , we adopt this value as a lower bound. Here, we evaluate the average gate fidelity of a π gate as a representative example of single-qubit gate fidelity. The qubit gate operation can be performed faster than the timescale of quasistatic noise and can be corrected through feedback operations [37–42]. Therefore, we first consider only high-frequency noise here. The average fidelity over all possible input states can be calculated as [36]

$$F_1(\delta) = \frac{1}{6} \left[3 + e^{-2t_g \gamma'_s} + 2e^{-t_g \gamma'_s} \cos(t_g \delta) \right], \quad (5)$$

where $t_g = 1/(2f_r^s)$ is the gate duration, and δ is the detuning between the spin resonance frequency and the frequency of the MW to the Co/Ti/(Nb-Ti-N) electrode. F_1 reaches its maximum when Λ^2 matches $\gamma_s/(\gamma_c - \gamma_s)$, where the ratio f_r^s/γ'_s is maximized. When quasistatic noise is also taken into account, with δ following a Gaussian distribution with standard deviation $\sqrt{2}\gamma_s^{*'}$, the average fidelity becomes

$$\bar{F}_1 = \frac{1}{6} \left[3 + e^{-2t_g \gamma'_s} + 2e^{-t_g \gamma'_s} e^{-\frac{(t_g \gamma_s^{*'})^2}{2}} \right]. \quad (6)$$

In Fig. 4(a), we show $1 - F_1(0)$ for scenarios 1, 2, and 3 and $1 - \bar{F}_1$ for scenario 1 as a function of Λ . For all the three scenarios, since the quasistatic noise is not significantly larger than the high-frequency noise, including it does not lead to a noticeable degradation in gate fidelity.

Next, we consider a two-qubit gate. Two-qubit operations can be performed on two electrons coupled to a resonator with the same geometry as the one used in this work [13]. Moreover, by unfolding the resonator, we can implement remote two-qubit gates with electrons positioned at each end of the resonator, as demonstrated in Refs. [6–10]. Among the various remote two-qubit gates [43], we focus on the i SWAP gate, which has been both theoretically analyzed [17,44] and experimentally realized for semiconductor spin qubits [10]. For this gate, spins are

TABLE II. The charge and spin loss rates assumed in three different scenarios. The charge loss rate values are experimentally measured at the charge sweet spot ($\epsilon = 0$), as reported in Refs. [12,13]. The spin loss rates are theoretical estimates from Ref. [15], where three scenarios are considered, each dominated by a different source of spin decoherence: thermal magnetic noise, hyperfine interaction with nuclear spins in natural neon ($^{\text{nat}}\text{Ne}$), or hyperfine interaction in isotopically enriched ^{22}Ne containing 1 ppm (part per million) of ^{21}Ne . For γ_s , as the relaxation time $T_{s,1}$ is expected to be long [15], $\gamma_s \approx \gamma_{s,\phi}$. For γ_c , the measurements reported in Refs. [12,13] show that $\gamma_c \approx \gamma_{c,1}/2$.

	Scenario 1 (thermal)	Scenario 2 ($^{\text{nat}}\text{Ne}$)	Scenario 3 (^{22}Ne)
$\gamma_c/2\pi$	$\frac{1}{2T_{c,1}} = \frac{1}{2 \times 11.6 \mu\text{s}} = 43.1 \text{ kHz}$ (Experimental value reported in Ref. [13])	$\frac{1}{2T_{c,1}} = \frac{1}{2 \times 48.2 \mu\text{s}} = 10.4 \text{ kHz}$ (Experimental value reported in Ref. [12])	
γ_c^*	$\frac{1}{T_{c,2}^*} = \frac{1}{8.2 \mu\text{s}} = 122 \text{ kHz}$ (Experimental value reported in Ref. [13])	$\frac{1}{T_{c,2}^*} = \frac{1}{42.8 \mu\text{s}} = 23.4 \text{ kHz}$ (Experimental value reported in Ref. [12])	
$\gamma_s/2\pi$	$\frac{1}{T_{s,\phi}} = \frac{1}{0.17 \text{ ms}} = 5.88 \text{ kHz}$ (Calculated value for thermal noise [15])	$\frac{1}{T_{s,\phi}} = \frac{1}{30 \text{ ms}} = 33.3 \text{ Hz}$ (Calculated value for $^{\text{nat}}\text{Ne}$ [15])	$\frac{1}{T_{s,\phi}} = \frac{1}{81 \text{ s}} = 0.012 \text{ Hz}$ (Calculated value for ^{22}Ne [15])
γ_s^*	$\frac{1}{T_{s,2}^*} = \frac{1}{0.16 \text{ ms}} = 6.25 \text{ kHz}$ (Calculated value for $^{\text{nat}}\text{Ne}$ [15])		$\frac{1}{T_{s,2}^*} = \frac{1}{0.43 \text{ s}} = 2.33 \text{ Hz}$ (Calculated value for ^{22}Ne [15])

coupled dispersively to the resonator via virtual photons, with the frequencies of the two qubits significantly detuned from that of the resonator ($\Delta_s = \beta g_s$, where $\beta \gg 1$). In the dispersive regime, we can largely suppress the effect of photon decay. The fidelity of the *i*SWAP gate is calculated as follows [17]:

$$F_2 = 1 - \frac{2\pi}{5g_s} \left(2\gamma_s'\beta + \frac{\kappa}{\beta} \right). \quad (7)$$

Since the quasistatic noise is not substantially greater than the high-frequency noise, we assume it to be negligible during two-qubit gate operations and take into account only the high-frequency noise. F_2 reaches its maximum when Λ^2 matches $(\gamma_s + \kappa/2\beta^2)/(\gamma_c - \gamma_s)$. In Fig. 4(b), we show $1 - F_2$ as a function of Λ for scenarios 1, 2, and 3. For scenario 1, where γ_s is on the same order as γ_c but still smaller, and satisfies $\gamma_s \gg \kappa/2\beta^2$, F_2 reaches

its maximum when $\Lambda^2 = \gamma_s/\gamma_c$ [vertical dashed green line in Fig. 4(b)]. For scenarios 2 and 3, where $\gamma_s \ll \gamma_c, \kappa/2\beta^2$, F_2 reaches its maximum when $\Lambda^2 = \kappa/(\beta^2 2\gamma_c)$ [vertical dashed blue line in Fig. 4(b)]. At this point, $F_2 = 1 - (4\pi/5)\sqrt{2\gamma_c\kappa/g_s^2}$, which is independent of both the spin-photon coupling strength and the spin decoherence rate, and reduces to the same form as that derived for semiconductor spin qubits in Ref. [17, Eq. (7)]. To evaluate the performance of the system under experimentally feasible conditions in the dispersive regime, we consider the following set of parameters as an example: $\omega_r/2\pi = 4.8 \text{ GHz}$, $\epsilon = 0$, $2t_c/2\pi = 8.0 \text{ GHz}$, $b_{\perp}/2\pi = 1 \text{ GHz}$, and $b_{\parallel}/2\pi = 5.1 \text{ GHz}$. Under these conditions, we obtain $\Lambda = 0.2$ and $\beta = 10$. As shown in Fig. 4(b), at $\Lambda = 0.2$, the single-qubit gate fidelity reaches 99.877% for scenario 1, 99.992% for scenario 2, and 99.996% for scenario 3, while the *i*SWAP gate fidelity reaches 99.34% for scenario 1 and 99.92% for scenarios 2 and 3.

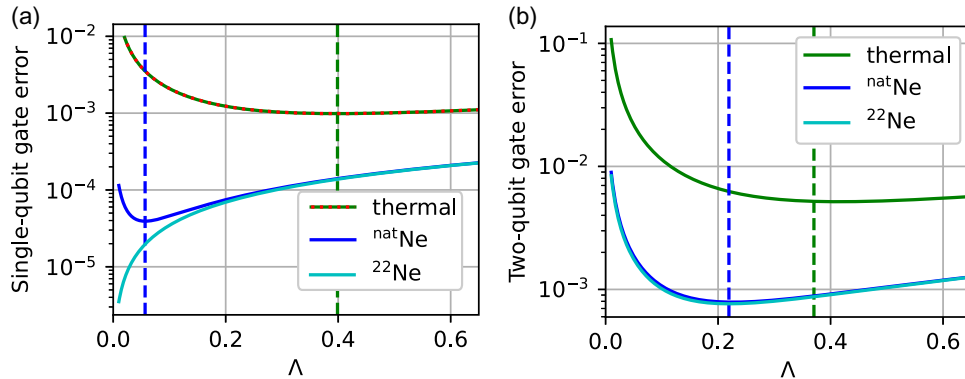


FIG. 4. (a) The single-qubit gate error as a function of the spin-charge coupling strength Λ . For scenario 1, the dotted green and red lines show $1 - F_1(0)$ and $1 - \bar{F}_1$, with the minimum at $\Lambda = 0.40$ (vertical dashed green line). For scenarios 2 and 3, the blue and cyan lines show $1 - F_1(0)$, respectively. For scenario 2, the error is minimized at $\Lambda = 0.057$, as indicated by the vertical blue dashed line. (b) The *i*SWAP gate error $1 - F_2$ as a function of Λ , calculated from Eq. (7), for scenario 1 (green line) with a minimum at $\Lambda = 0.37$ (vertical dashed green line), and for scenarios 2 and 3 (blue and cyan lines) with a minimum at $\Lambda = 0.22$ (vertical dashed blue line). The parameter $\beta = \Delta_s/g_s$ is set to 10 and $\kappa/2\pi = 0.1 \text{ MHz}$ in all cases.

The gate fidelities might be further improved by dynamically tuning $2t_c$ (or ϵ) via voltage gate pulses to optimize Λ independently for single- and two-qubit gate operations. For instance, in scenario 3, with isotopically purified neon, reducing Λ can further enhance the single-qubit gate fidelity to exceed 99.999%. This can be achieved by increasing $2t_c$; however, tuning to high $2t_c$ may be challenging if it is determined by surface roughness, as discussed in Sec. III. Alternatively, reducing $b_\perp$ through device design also lowers Λ , but this comes at the cost of a reduced spin Rabi frequency f_r^s , which increases the gate time t_g and ultimately degrades fidelity. Moreover, at this level of high gate fidelity, instrumental factors such as phase noise in the control electronics become increasingly important. To fully benefit from isotopic purification, it is therefore crucial to also consider these instrumental limitations [45]. Since the *i*SWAP gate fidelity is limited not by the spin loss rate but by the charge and resonator loss rates, isotopic purification of neon does not yield a significant improvement in the two-qubit gate performance, at least for the scheme considered here.

VII. CONCLUSIONS

In conclusion, we experimentally demonstrated neon and electron depositions using Nb-Ti-N nanowire resonators, aiming for the realization of spin-photon coupling and high-fidelity spin-qubit gates in the future. The presence of neon and electrons was confirmed by a decrease in the resonance frequency without compromising the quality factor of the resonator, supporting the suitability of Nb-Ti-N resonators for future qubit measurements. A closer examination of the response of the resonator to electron deposition suggests that the Drude model fails to capture the behavior of the system, likely due to electron localization induced by surface roughness in the neon. At present, the control over thickness and surface roughness in the neon deposition process is not well established and requires improvement. Additionally, we investigated the configuration of cobalt magnets needed for realizing spin-photon coupling and spin-qubit gates. By combining the enhanced charge-photon coupling enabled by the high impedance of the Nb-Ti-N resonator with the calculated magnetic field gradient produced by the cobalt magnets, we estimated that spin-photon coupling can reach the strong-coupling regime, allowing for the implementation of high-fidelity spin-qubit gates.

Future work will involve the experimental implementation of integrated micromagnets and additional dc electrodes, as theoretically proposed in this work. To avoid degradation of the quality factor of the resonator in such implementations, mitigation strategies such as perforated ground planes [24] and dc filters [46,47] will be necessary. In this context, the high quality factor demonstrated

here provides a target and benchmark for evaluating the effectiveness of these approaches.

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DATA AVAILABILITY

The data that support the findings of this paper are openly available [48].

APPENDIX A: SAMPLE FABRICATION

Samples are fabricated on high-resistivity ($\rho \geq 10 \text{ k}\Omega \text{ cm}$) (100)-oriented intrinsic silicon substrates. The substrates are cleaned in a piranha solution for 15 min and then rinsed with deionized (DI) water. To remove the surface oxide, they are dipped in buffered hydrofluoric acid (BHF) for 5 min, followed by thorough rinsing with DI water. A 20-nm-thick Nb-Ti-N thin film is prepared by reactive magnetron sputtering using an 8-in. alloy target with a weight ratio of Nb:Ti = 5:1. During deposition, the substrate is kept at ambient temperature, with a dc sputtering current of 3 A and a gas pressure of 2 mTorr (Ar:N₂ = 100:37). After cutting the wafer with the Nb-Ti-N thin film into $10 \times 10 \text{ mm}$ pieces, it undergoes oxygen plasma cleaning for 1 min to remove organic contamination, followed by a 1-min BHF clean to eliminate any remaining residues. A 100-nm-thick layer of ZEP520A resist, diluted in anisole, is spin coated onto the substrate. The coated substrate is then baked at 180°C for 3 min. Electron-beam lithography (EBL) is carried out using an ELS-F125 system with an electron dose of $230 \mu\text{C}/\text{cm}^2$ and a beam current of 1 nA to define the microwave circuitry. After exposure, the resist is developed in ZED-N50 for 2 min, followed by a 30-s rinse in IPA. The pattern is transferred to the film by reactive ion etching (RIE) using a gas mixture of O₂ for 25 s and CF₄ at 50 sccm, 100 W, and 10 Pa for 75 s. This process results in the successful fabrication of the Nb-Ti-N resonator on the chip, with the underlying 60-nm Si layer also being etched.

APPENDIX B: SIMULATION FOR NEON AND ELECTRON DEPOSITION

In Figs. 5(a) and 5(b), we represent resonant peaks measured for resonators 1 and 2 with and without neon and electrons at 3.1 K and around 7 mK, and in Fig. 5(c) we show the resonance-frequency shifts calculated as a function of neon thickness using COMSOL simulations. The neon thickness is defined from the bottom of the etched Si substrate. For the “three-dimensional (3D) rf model” described in Appendix B 1, since the resonator is defined on a 2D plane located 80 nm above the bottom of the etched Si substrate, the graph starts from a neon thickness of 80 nm.

The resonance-frequency shift measured around 7 mK can be attributed to both the effects of neon and electron deposition. By comparing the experimental and simulation results, we can thus set an upper bound for the neon thickness.

To obtain a reliable estimate of the thickness, we employed and compared two independent simulation methods: a 3D electromagnetic simulation using COMSOL (see Appendix B 1) and a simplified 2D transmission line calculator (TLC; see Appendix B 2). The 3D model captures detailed geometric and field-distribution effects, while the TLC approach provides fast intuitive access to the dependence on dielectric loading. Consistent results

from both methods increase our confidence in the estimated neon thickness.

Thus, the neon thickness deposited on resonator 1 was estimated to be 160 nm in both models. For resonator 2, the resonance frequency is closer to the simulation result for a width of 150 nm than to that for 200 nm (see Table I). Therefore, using the simulation result from the TLC (see Appendix B 2) for a width of 150 nm, the neon thickness was estimated to be 270 nm.

1. Three-dimensional (3D) rf model

The simulation was performed using COMSOL Multiphysics v. 6.2. We began the simulation by creating a 3D model with the “Radio Frequency” module loaded for the physics setup, within which the “Electromagnetic Waves, Frequency Domain (emw)” interface is used. The study type was set to Frequency Domain and the MUMPS direct solver was used to compute the response of the system subjected to harmonic excitation over a range of frequencies.

We use a 3D model with dimensions of $2000\ \mu\text{m} \times 2000\ \mu\text{m} \times 800\ \mu\text{m}$, where the center of the resonator was roughly positioned at the origin. For simulating the resonator response in the absence of neon, the region $301.3 < z < 800\ \mu\text{m}$ was defined as vacuum and the region $0 <$

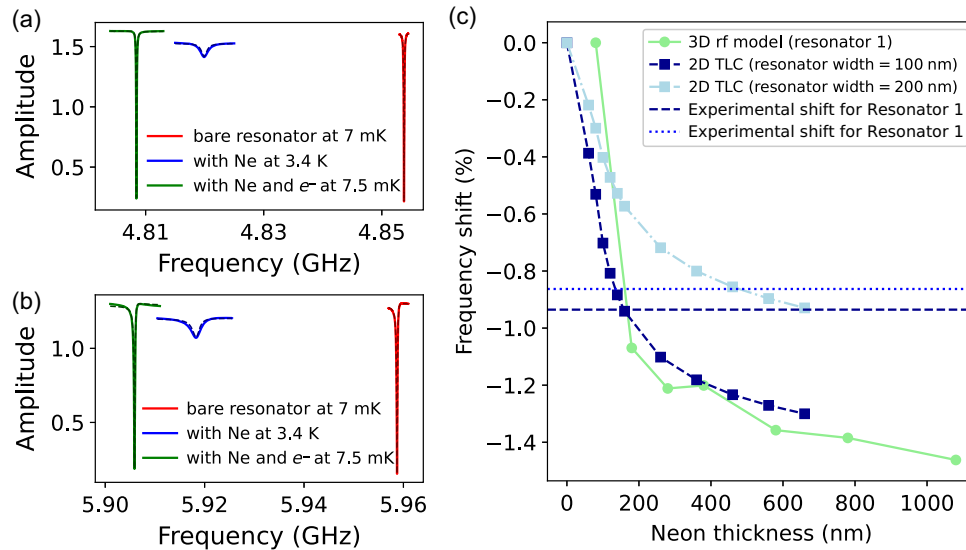


FIG. 5. (a) The resonance peaks measured for the bare resonator 1 at 7 mK (red), with neon at 3.4 K (blue), and with neon and electrons at 7.5 mK (green). The corresponding internal quality factors (Q_{int} : 2.34×10^5 , 4.66×10^3 , and 2.22×10^5), external quality factors (Q_{ext} : 3.19×10^4 , 4.33×10^3 , and 3.36×10^4), and resonance frequencies (f_r : 4.854 GHz, 4.82 GHz, and 4.808 GHz) are given for red, blue, and green, respectively. (b) The resonance peaks measured for the bare resonator 2 at 7 mK (red), with neon at 3.4 K (blue), and with neon and electrons at 7.5 mK (green). The corresponding internal quality factors (Q_{int} : 1.91×10^5 , 3.9×10^3 , and 1.42×10^5), external quality factors (Q_{ext} : 2.16×10^4 , 3.49×10^3 , and 1.97×10^4), and resonance frequencies (f_r : 5.959 GHz, 5.918 GHz, and 5.906 GHz) are given for red, blue, and green, respectively. (a),(b) The dotted lines represent the fits. (c) The resonance-frequency shift obtained from the COMSOL simulation as a function of the neon thickness. The vertical dotted and dashed lines represent the maximum measured resonance-frequency shifts of resonators 1 and 2 caused by the neon and electron deposition, -0.94% and -0.86% , respectively. The resonance-frequency shifts for resonator 1 are simulated using the three-dimensional (3D) rf model.

$z < 300 \mu\text{m}$ was defined as silicon (the relative permittivity of silicon was set to 11.4). Thirteen layers with thickness of 100 nm were added from $300 < z < 301.3 \mu\text{m}$, each of which could be set as neon ($\epsilon_r = 1.244$) or vacuum to simulate a neon layer of various thicknesses.

We set a 2D plane at $z = 300 \mu\text{m}$, where the resonators, the feedline, and the grounds are defined. The transmission of the electromagnetic field is solved using lumped ports 1 and 2, which are defined between each end of the feedline and the ground plane, each with a characteristic impedance of 50Ω . The transmission parameter S_{21} is then plotted as a function of the frequency f of the microwave signal sent through the feedline.

We account for the kinetic inductance of Nb-Ti-N by modeling it as a complex conductivity within the “Transition Boundary Condition” applied to the resonator. According to Ref. [49], the current density $\vec{J}$ can be expressed as

$$\vec{J} = \sigma \vec{E},$$

where $\vec{E}$ is the electric field, and the conductivity σ for a thin film is given by

$$\sigma = \frac{1}{i\mu_0\omega\lambda^2}, \quad (\text{B1})$$

where $\omega = 2\pi f$, μ_0 is the magnetic permeability, and λ is the magnetic penetration depth. In the COMSOL simulation, the electrical conductivity for the “Transition Boundary Condition” of the resonator is set according to Eq. (B1). When the sample thickness is comparable to or smaller than the skin depth, as is the case here, we deselect the “Electrically thick layer” option and specify a layer thickness of $D = 20 \text{ nm}$.

The feedline and grounds are modeled as “Perfect Electric Conductor,” even though they are also made of Nb-Ti-N, because their kinetic inductance is expected to have a negligible impact on the resonator response for this geometry. The quality factor of the resonator computed in this COMSOL simulation is determined solely by the external quality factor, as we treat silicon as a lossless dielectric in the simulation.

In Table I, we show the comparison between the measured data and the COMSOL simulation with $\lambda = 390 \text{ nm}$. Compared to resonator 1, resonators 2 and 3 exhibit deviations from the simulation results based on the intended width and gap dimensions. This is believed to be due to the width of resonators 2 and 3 being smaller than intended, resulting in lower frequencies compared to the simulation.

2. “Transmission Line Calculator” COMSOL model

We modeled the nanowire resonator as a coplanar waveguide using the “Transmission Line Calculator” COMSOL model. Although the actual resonator is not a straight

nanowire, in this model, it is approximated as a transmission line aligned with the x' axis (designating this extending direction as the x' axis). Electromagnetic fields, specifically transverse electromagnetic (TEM) waves, propagate along the x' axis. Here, we also assume that the x' components of the electric and magnetic fields are small and the propagating mode is deduced from separate magnetic and electric analyses.

The width of the resonator is defined along the y' axis as 150 nm and the thickness along the z axis as 20 nm, and it extends indefinitely along the x' axis. In this model, we define metal, silicon, neon, and electrons in the y' - z plane.

a. Simulation of neon thickness

To estimate the neon thickness, the neon thickness is considered in the z direction, growing sequentially. We obtained the shunt capacitance per unit length, denoted as C_l . The frequency shift due to the presence of neon was then calculated using the formula $\Delta f = \sqrt{C_l^{\text{without Ne}}/C_l^{\text{with Ne}}} - 1$, where $C_l^{\text{without Ne}}$ and $C_l^{\text{with Ne}}$ represent the shunt capacitance per unit length without and with neon, respectively.

b. Simulation with the presence of electrons: Drude model

To incorporate the effect of electrons, we first use the Drude model [50] based on the same 2D framework described above. The neon thickness is 270 nm and the surface electrons are assumed to be located 2.5 nm above the neon surface, spreading along the x' - y' plane. The equation of motion for an electron is given by

$$m_e \left(\frac{\partial \mathbf{v}}{\partial t} + \frac{\mathbf{v}}{\tau} \right) = -e\mathbf{E}, \quad (\text{B2})$$

where $\mathbf{E} = \mathbf{E}_0 \exp(i\omega t)$ is the electric field acting on the electron [51] and $\mathbf{v}$ is the electron velocity. Here, τ is the scattering time and m_e is the electron mass. The solution to Eq. (B2) is

$$\mathbf{v} = -\frac{e\tau}{m_e} \frac{1}{1 + i\omega\tau} \mathbf{E}_0 \exp(i\omega t). \quad (\text{B3})$$

The corresponding 2D current density is $\mathbf{j} = -en_e\mathbf{v} = \sigma^{2D}\mathbf{E}_0 \exp(i\omega t)$, leading to the well-known Drude model:

$$\sigma^{2D} = \frac{e^2 n_e \tau}{m_e} \frac{1}{1 + i\omega\tau}. \quad (\text{B4})$$

To integrate this into the COMSOL model, the electron conductivity is defined as

$$\sigma_{ij} = \frac{1}{\delta z} \frac{e^2 n_e \tau}{m_e} \frac{1}{1 + i\omega\tau} \quad (\text{B5})$$

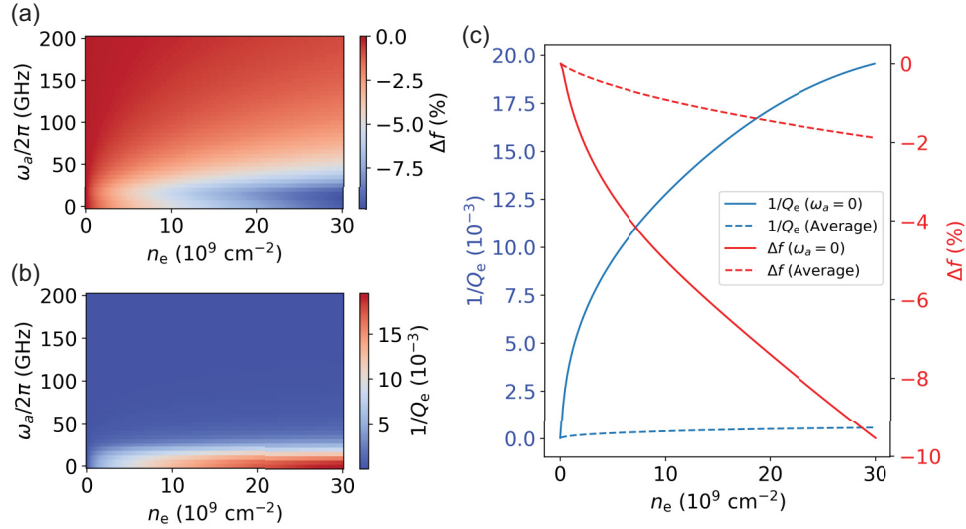


FIG. 6. (a),(b) The (a) resonance-frequency shift Δf and (b) inverse quality factor $1/Q_e$, due to the presence of electrons, obtained from simulations using Eq. (B10) as a function of the electron density n_e and the confinement frequency $\omega_a/2\pi$. (c) The simulated $1/Q_e$ (blue line) and Δf (red line) as functions of n_e for $\omega_a/2\pi = 0$ GHz (solid line) and the average over $\omega_a/2\pi \in [0, 200]$ GHz with weights $\propto \exp(\hbar\omega_a/k_B T)$ (dashed line).

for $i = j = x'$ or $i = j = y'$, and $\sigma_{ij} = 0$ otherwise, where $\delta z = 1$ nm is the arbitrarily defined thickness of the area considered as the electron layer. To check the validity of this choice, simulations were also conducted with $\delta z = 2.5$ nm and it was confirmed that the differences in the results were negligibly small. Note that as the x' component of the electric and magnetic fields is small, $\sigma_{y'y'}$ mainly contributes to the resonance peak change. We obtained the shunt capacitance per unit length and calculated the frequency shift as $\Delta f = \sqrt{C_l^{\text{without } e^-} / C_l^{\text{with } e^-}} - 1$, where $C_l^{\text{without } e^-}$ and $C_l^{\text{with } e^-}$ represent the shunt capacitance per unit length without and with electrons, respectively. The attenuation constant α was extracted from the complex propagation constant $\gamma = \alpha + i\beta$, computed in the simulation. The internal quality factor due to the electrons was then obtained via [52]

$$Q_e = \frac{\pi}{2\alpha l}. \quad (\text{B6})$$

Using this approach, we simulated the frequency shift Δf and electron-induced internal loss $1/Q_e$ with $\tau = 1.9$ ps, corresponding to the $\omega_a = 0$ case in Fig. 6. As discussed in the main text, the resulting $1/Q_e$ from the simulation is nearly an order of magnitude larger than the experimentally measured value, which remains around 4×10^{-4} at 3.4 K and shows no discernible dependence on the electron density [see Fig. 2(c)].

c. Simulation with the presence of electrons: Lorentz model

To explain the discrepancy between the measured and simulated internal quality factors, we consider electron localization induced by surface roughness. As one possible modeling approach, we assume that all electrons are confined in harmonic potential wells characterized by a frequency ω_a . The electron dynamics are then governed by the equation of motion:

$$m_e \left(\frac{\partial \mathbf{v}}{\partial t} + \frac{\mathbf{v}}{\tau} \right) = -e\mathbf{E} - m_e \omega_a^2 \mathbf{x}, \quad (\text{B7})$$

where $\mathbf{x}$ is the position of the electron, taking the origin as the center of the harmonic potential that it experiences, with a confinement frequency ω_a . Here, τ is the scattering time and m_e is the electron mass. The solution to Eq. (B7) is

$$\mathbf{v} = i\omega \mathbf{x} = -\frac{e\tau}{m_e} \frac{1}{1 + i\left(\omega - \frac{\omega_a^2}{\omega}\right)\tau} \mathbf{E}_0 \exp(i\omega t). \quad (\text{B8})$$

The resulting modified surface conductivity is given by

$$\sigma^{2D} = \frac{e^2 n_e \tau}{m_e} \frac{1}{1 + i\left(\omega - \frac{\omega_a^2}{\omega}\right)\tau}. \quad (\text{B9})$$

To integrate this into the COMSOL model, the electron conductivity is now redefined as

$$\sigma_{ij} = \frac{1}{\delta z} \frac{e^2 n_e \tau}{m_e} \frac{1}{1 + i\left(\omega - \frac{\omega_a^2}{\omega}\right)\tau}. \quad (\text{B10})$$

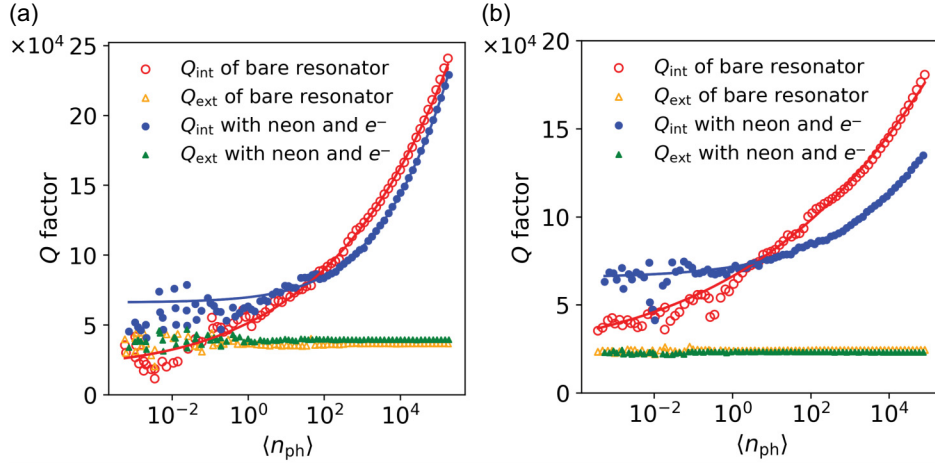


FIG. 7. The measured quality-factor power dependence for (a) resonator 1 and (b) resonator 2 as a function of the number of photons in the resonators, $\langle n_{ph} \rangle$. The open red circles and solid blue circles represent Q_{int} of the bare resonator and the resonator with electrons and neon, respectively. The open orange triangles and solid green triangles represent Q_{ext} of the bare resonator and the resonator with electrons and neon, and we found that they are independent of $\langle n_{ph} \rangle$, with no observable changes before and after the deposition of electrons and neon. The solid red and blue lines are fits to Q_{int} using Eq. (C1). (a) For resonator 1, $Q_{TLS,0}/F$ increases from $(1.94 \pm 0.87) \times 10^4$ to $(6.64 \pm 0.92) \times 10^4$, n_{sat} increases from $(1 \pm 5) \times 10^{-3}$ to $(3.0 \pm 1.2) \times 10^2$, and β increases from 0.265 ± 0.051 to 0.377 ± 0.063 after the deposition of neon and electrons. Estimating Q_{other} from the fit is difficult due to the lack of data at high photon numbers. However, we find that Q_{other} changes from $(2 \pm 49) \times 10^7$ to $(2 \pm 37) \times 10^7$ after the deposition of neon and electrons. (b) For resonator 2, $Q_{TLS,0}/F$ increases from $(2.2 \pm 2.2) \times 10^4$ to $(6.8 \pm 1.6) \times 10^4$, n_{sat} increases from $(1 \pm 14) \times 10^{-5}$ to $(4.7 \pm 2.8) \times 10^2$, and β increases from 0.190 ± 0.060 to 0.256 ± 0.059 after the deposition of neon and electrons. Additionally, Q_{other} changes from $(2 \pm 10) \times 10^6$ to $(2 \pm 6) \times 10^6$ after the deposition of neon and electrons. For both resonators, the errors are estimated at the 95% confidence level. The large uncertainties in n_{sat} and Q_{other} arise because we did not reach a high enough power to fully saturate the two-level-system (TLS). However, the values of n_{sat} and Q_{other} did not affect the fitting results of the other parameters.

Using the same approach as described above, we simulated the frequency shift Δf and electron-induced internal loss $1/Q_e$ with $\tau = 1.9$ ps, for trap frequencies $\omega_a/2\pi$ ranging from 0 to 200 GHz [see Figs. 6(a) and 6(b)]. Introducing a finite ω_a enhances the simulated resonance-frequency shift around $\omega_a \sim \sqrt{\omega/\tau}$, where the imaginary part of σ^{2D} peaks. Meanwhile, the degradation of the simulated Q_{int} is reduced as ω_a increases, since the real part of σ^{2D} decreases monotonically with increasing ω_a .

To explain the experimentally observed absence of degradation in Q_{int} , we take the following approach. We assume that the number of traps is significantly larger than the number of electrons and that the trap depths ω_a are uniformly distributed. The thermal energy at $T = 3.4$ K is $k_B T \approx 70$ GHz $\times h$. Electrons trapped in deeper potential wells ($\hbar\omega_a \gg k_B T$) remain trapped longer, whereas those in shallower traps escape more quickly. Following previous studies on thermal escape time from quantum wells [53,54], we assume that the thermal escape time of an electron from a trap scales as $\exp(\hbar\omega_a/k_B T)$ and, similarly, the probability of an electron occupying a trap at a given ω_a follows the same dependence. We calculate and plot Δf and Q_e by integrating over trap depths ranging from $\omega_a/2\pi = 0$ to 200 GHz. As shown in Fig. 6(c), this model successfully reproduces the experimentally

observed behavior where Q_{int} does not degrade at $\Delta f = -0.9\%$, yielding $1/Q_e = 3.7 \times 10^{-4}$. Note that the effective electron density n_e in this model refers only to electrons trapped within this frequency range, estimated to be approximately $9.6 \times 10^9 \text{ cm}^{-2}$ for $\Delta f = -0.9\%$. Electrons trapped deeper than 200 GHz do not significantly contribute to Δf or Q_{int} within this framework, as seen in Figs. 6(a) and 6(b).

APPENDIX C: POWER DEPENDENCE OF Q_{int}

In Fig. 7, we show Q_{int} and Q_{ext} of resonators 1 and 2 as a function of the average photon number $\langle n_{ph} \rangle$ in the resonator, measured for the bare resonator, and with neon and electrons present. Since this measurement was performed at 10 mK, we assume that the effect of electrons is minimal under our surface conditions. As shown in Fig. 2, the quality factor does not degrade upon electron loading. Thus, the observed variations in Q_{int} are primarily associated with the presence or absence of neon. The solid lines in Fig. 7 represent fitting curves based on the following TLS loss model:

$$\frac{1}{Q_{int}} = \frac{F}{Q_{TLS,0} \sqrt{1 + \left(\frac{\langle n_{ph} \rangle}{n_{sat}}\right)^\beta}} + \frac{1}{Q_{other}}, \quad (C1)$$

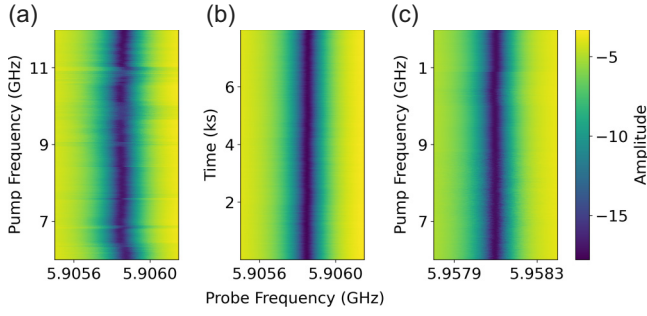


FIG. 8. Resonance peaks of resonator 2 measured at 10 mK. We observed that the resonance frequency changed only when both neon and electrons were present. (a) Two-tone measurement with both neon and electrons. (b) Stability measurement of the resonance peak over time. (c) Two-tone measurement following the removal of electrons and neon.

where F is the participation ratio of the electric energy in the region in which the TLS exists, $1/Q_{\text{TLS},0}$ is the intrinsic TLS loss in the zero-photon and zero-temperature limit, n_{sat} is the saturation photon number, and β is an empirical parameter that accounts for TLS dynamics. Since the data points with lower $\langle n_{\text{ph}} \rangle$ have a lower signal-to-noise ratio (SNR), we set the fitting weight w_i of each data point i proportional to its SNR. The fitting results demonstrate a good fit and suggest that Q_{int} of the Nb-Ti-N resonators is limited by TLS loss [55–58]. With neon and electrons, $Q_{\text{TLS},0}/F$ increases by a factor of approximately 3 compared to the bare resonator. This is likely due to the gap between the resonator and the ground being filled with neon, which has a higher dielectric constant than vacuum [59]. As a result, the participation ratio F decreases [55]. Additionally, n_{sat} also increases, indicating that more microwave power is required to saturate the TLS due to the presence of neon. However, the observed changes are larger than expected based on the relative permittivity of neon, which is 1.244. This suggests that some other mechanism might also be involved, and further study is needed.

APPENDIX D: TWO-TONE MEASUREMENT

In Fig. 8, we show the two-tone measurement on resonator 2. At room temperature, a probe signal with a frequency near the resonance frequency is combined with a pump signal in the 6–12 GHz range and injected into port 1. The pump signal excites from the ground orbital state to the orbital excited state when their energy difference corresponds to the pump signal frequency, causing a shift in the resonance frequency.

APPENDIX E: MICROMAGNET SIMULATION

The numerical calculation of the local magnetic fields in different directions is done by modeling each Co micromagnet as a uniformly magnetized rectangular block. We treat the block as a sum of many small magnetic dipoles and integrate their fields over the length, width, and thickness of the block [34].

In this work, we evaluate two Co blocks with magnetization $M = 1.7$ T and dimensions $1.5 \mu\text{m} \times 1.5 \mu\text{m} \times t$, where the t represents the different thickness to be simulated. The magnetic field along the axis of interest was calculated with the MAGPYLIB PYTHON software package.

As shown in Fig. 3(b), dB_z/dy is maximized for $\Delta z \approx 146$ nm. We also simulate the displacement sensitivity. A 10-nm shift of the micromagnets along y changes dB_z/dy only slightly [dashed orange line in Fig. 9(a)]; this indicates that small misalignments have little effect. A larger 100-nm shift (dotted green line) increases dB_z/dy at the same Δz . Thus, within normal fabrication tolerances, the gradient—and therefore the coupling strength—can be maintained or even improved as the electron moves closer to one of the magnets.

The absolute magnetic fields with misalignment along the y direction were also simulated [see Figs. 9(b) and 9(c)]. The fields at the resonator position stay well below the Nb-Ti-N critical field. The field at the resonator position remains well below the critical field of Nb-Ti-N. Therefore, the micromagnets are not expected to degrade the quality factor of the resonator. In addition, we note that

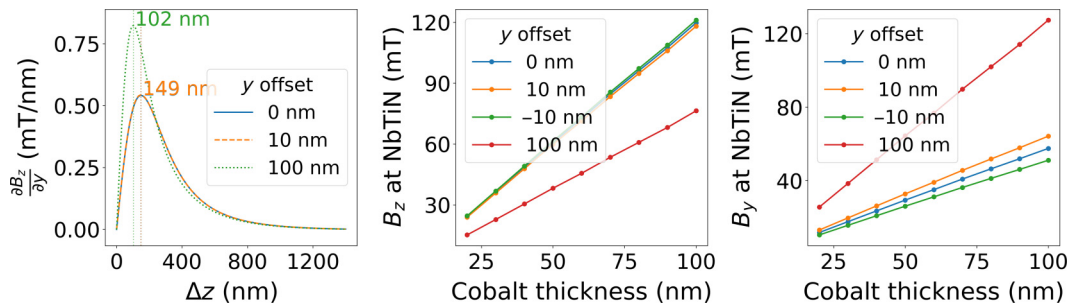


FIG. 9. (a) $\partial B_z / \partial y$ as a function of Δz for a 100-nm-thick cobalt magnet with misalignment along the y direction. (b) B_z and (c) B_y at the resonator position as functions of cobalt thickness, both with misalignment along the y direction.

displacement along the x direction has a negligible impact on both the magnetic field and its gradient.

APPENDIX F: KINETIC INDUCTANCE

From separate measurements at room temperature, we estimated the magnetic penetration depth at zero temperature for the Nb-Ti-N film with a thickness of $D = 20$ nm, used in the experiments, to be $\lambda = 390$ nm. The inductance per square $L_{\square}$ is calculated as $L_{\square} = \mu_0 \lambda^2 / D = 9.6$ pH. The total inductance of the resonator is determined by multiplying $L_{\square}$ by the number of squares that fit within the geometry. Thus, the kinetic inductance of resonator 1 is calculated as $L = L_{\square} l / w = 9.6 \text{ pH} \times 1.45 \text{ mm} / 100 \text{ nm} = 139 \text{ nH}$, where l and w are the length and the width of the resonator, respectively. As shown in Table I, the resonance frequencies obtained from both simulation and experiment are in good agreement for resonator 1, thereby validating this inductance estimation.

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- [1] K. D. Petersson, L. W. McFaul, M. D. Schroer, M. Jung, J. M. Taylor, A. A. Houck, and J. R. Petta, Circuit quantum electrodynamics with a spin qubit, *Nature* **490**, 380 (2012).
 - [2] X. Mi, J. V. Cady, D. M. Zajac, P. W. Deelman, and J. R. Petta, Strong coupling of a single electron in silicon to a microwave photon, *Science* **355**, 156 (2017).
 - [3] A. Stockklauser, P. Scarlino, J. V. Koski, S. Gasparinetti, C. K. Andersen, C. Reichl, W. Wegscheider, T. Ihn, K. Ensslin, and A. Wallraff, Strong coupling cavity qed with gate-defined double quantum dots enabled by a high impedance resonator, *Phys. Rev. X* **7**, 011030 (2017).
 - [4] A. J. Landig, J. V. Koski, P. Scarlino, U. C. Mendes, A. Blais, C. Reichl, W. Wegscheider, A. Wallraff, K. Ensslin, and T. Ihn, Coherent spin-photon coupling using a resonant exchange qubit, *Nature* **560**, 179 (2018).
 - [5] N. Samkharadze, A. Bruno, P. Scarlino, G. Zheng, D. P. DiVincenzo, L. DiCarlo, and L. M. K. Vandersypen, High-kinetic-inductance superconducting nanowire resonators for circuit qed in a magnetic field, *Phys. Rev. Appl.* **5**, 044004 (2016).
 - [6] X. Mi, M. Benito, S. Putz, D. M. Zajac, J. M. Taylor, G. Burkard, and J. R. Petta, A coherent spin-photon interface in silicon, *Nature* **555**, 599 (2018).
 - [7] F. Borjans, X. G. Croot, X. Mi, M. J. Gullans, and J. R. Petta, Resonant microwave-mediated interactions between distant electron spins, *Nature* **577**, 195 (2020).
 - [8] C. G. L. Böttcher, S. P. Harvey, S. Fallahi, G. C. Gardner, M. J. Manfra, U. Vool, S. D. Bartlett, and A. Yacoby, Parametric longitudinal coupling between a high-impedance superconducting resonator and a semiconductor quantum dot singlet-triplet spin qubit, *Nat. Commun.* **13**, 4773 (2022).
 - [9] P. Harvey-Collard, J. Dijkema, G. Zheng, A. Sammak, G. Scappucci, and Lieven M. K. Vandersypen, Coherent spin-spin coupling mediated by virtual microwave photons, *Phys. Rev. X* **12**, 021026 (2022).
 - [10] J. Dijkema, X. Xue, P. Harvey-Collard, M. Rimbach-Russ, S. L. de Snoo, G. Zheng, A. Sammak, G. Scappucci, and L. M. K. Vandersypen, Cavity-mediated iSWAP oscillations between distant spins, *Nat. Phys.* **21**, 168 (2025).
 - [11] X. Zhou, G. Koolstra, X. Zhang, G. Yang, X. Han, B. Dizdar, X. Li, R. Divan, W. Guo, K. W. Murch, D. I. Schuster, and D. Jin, Single electrons on solid neon as a solid-state qubit platform, *Nature* **605**, 46 (2022).
 - [12] X. Zhou, X. Li, Q. Chen, G. Koolstra, G. Yang, B. Dizdar, Y. Huang, C. S. Wang, X. Han, X. Zhang, D. I. Schuster, and D. Jin, Electron charge qubit with 0.1 ms coherence time, *Nat. Phys.* **20**, 116 (2024).
 - [13] X. Li, C. S. Wang, B. Dizdar, Y. Huang, Y. Wen, W. Guo, X. Zhang, X. Han, X. Zhou, and D. Jin, Noise-resilient host for electron qubit operation up to 0.4 K, *ArXiv:2502.01005*.
 - [14] X. Li, S. Zou, Q. Chen, and D. Jin, Electron charge coherence on a solid neon surface, *ArXiv:2507.20476*.
 - [15] Q. Chen, I. Martin, L. Jiang, and D. Jin, Electron spin coherence on a solid neon surface, *Quantum Sci. Technol.* **7**, 045016 (2022).
 - [16] M. Benito, X. Mi, J. M. Taylor, J. R. Petta, and G. Burkard, Input-output theory for spin-photon coupling in Si double quantum dots, *Phys. Rev. B* **96**, 235434 (2017).
 - [17] M. Benito, J. R. Petta, and G. Burkard, Optimized cavity-mediated dispersive two-qubit gates between spin qubits, *Phys. Rev. B* **100**, 081412(R) (2019).
 - [18] N. Samkharadze, G. Zheng, N. Kalhor, D. Brousse, A. Sammak, U. C. Mendes, A. Blais, G. Scappucci, and L. M. K. Vandersypen, Strong spin-photon coupling in silicon, *Science* **359**, 1123 (2018).
 - [19] A. G. Fowler, M. Mariantoni, J. M. Martinis, and A. N. Cleland, Surface codes: Towards practical large-scale quantum computation, *Phys. Rev. A* **86**, 032324 (2012).
 - [20] Y. Tokura, W. G. van der Wiel, T. Obata, and S. Tarucha, Coherent single electron spin control in a slanting Zeeman field, *Phys. Rev. Lett.* **96**, 047202 (2006).
 - [21] A. Jennings, X. Zhou, I. Grytsenko, and E. Kawakami, Quantum computing using floating electrons on cryogenic substrates: Potential and challenges, *Appl. Phys. Lett.* **124**, 120501 (2024).
 - [22] E. Kawakami, J. Chen, M. Benito, and D. Konstantinov, Blueprint for quantum computing using electrons on helium, *Phys. Rev. Appl.* **20**, 054022 (2023).
 - [23] D. I. Schuster, A. Fragner, M. I. Dykman, S. A. Lyon, and R. J. Schoelkopf, Proposal for manipulating and detecting spin and orbital states of trapped electrons on helium using cavity quantum electrodynamics, *Phys. Rev. Lett.* **105**, 040503 (2010).
 - [24] J. G. Kroll, F. Borsoi, K. L. van der Enden, W. Uilhoorn, D. de Jong, M. Quintero-Pérez, D. J. van Woerkom, A. Bruno, S. R. Plissard, D. Car, E. P. A. M. Bakkers, M. C. Cassidy, and L. P. Kouwenhoven, Magnetic-field-resilient superconducting coplanar-waveguide resonators for hybrid circuit quantum electrodynamics experiments, *Phys. Rev. Appl.* **11**, 064053 (2019).
 - [25] S. Probst, F. B. 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).
 - [26] S. Probst, Circle fit (2024), Available online: https://github.com/sebastianprobst/resonator_tools.

- [27] P. Leiderer, Surface electrons on solid quantum substrates: A brief review, *J. Low Temp. Phys.* **219**, 262 (2025).
- [28] A. D. Migone, J. G. Dash, M. Schick, and O. E. Vilches, Triple-point wetting of neon films, *Phys. Rev. B* **34**, 6322 (1986).
- [29] K. Kajita, A new two-dimensional electron system on the surface of solid neon, *Surf. Sci.* **142**, 86 (1984).
- [30] K. Kajita, Wigner crystallization of two dimensional electrons formed on the surface of solid neon, *J. Phys. Soc. Jpn.* **54**, 4092 (1985).
- [31] K. C. Nowack, F. H. L. Koppens, Y. V. Nazarov, and L. M. K. Vandersypen, Coherent control of a single electron spin with electric fields, *Science* **318**, 1430 (2007).
- [32] M. Pioro-Ladrière, T. Obata, Y. Tokura, Y.-S. Shin, T. Kubo, K. Yoshida, T. Taniyama, and S. Tarucha, Electrically driven single-electron spin resonance in a slanting Zeeman field, *Nat. Phys.* **4**, 776 (2008).
- [33] T. Kanai, D. Jin, and W. Guo, Single-electron qubits based on quantum ring states on solid neon surface, *Phys. Rev. Lett.* **132**, 250603 (2024).
- [34] J. R. Goldman, T. D. Ladd, F. Yamaguchi, and Y. Yamamoto, Magnet designs for a crystal-lattice quantum computer, *Appl. Phys. A* **71**, 11 (2000).
- [35] In analogy with Ref. [36], the decay due to quasistatic noise is assumed to follow a Gaussian form, $\exp(-(\gamma^*t)^2)$.
- [36] M. Benito, X. Croot, C. Adelsberger, S. Putz, X. Mi, J. R. Petta, and G. Burkard, Electric-field control and noise protection of the flopping-mode spin qubit, *Phys. Rev. B* **100**, 125430 (2019).
- [37] J. Kim, J. Yun, W. Jang, H. Jang, J. Park, Y. Song, M.-K. Cho, S. Sim, H. Sohn, H. Jung, V. Umansky, and D. Kim, Approaching ideal visibility in singlet-triplet qubit operations using energy-selective tunneling-based Hamiltonian estimation, *Phys. Rev. Lett.* **129**, 040501 (2022).
- [38] F. Berritta, T. Rasmussen, J. A. Krzywda, J. van der Heijden, F. Fedeale, S. Fallahi, G. C. Gardner, M. J. Manfra, E. van Nieuwenburg, J. Danon, A. Chatterjee, and F. Kuemmeth, Real-time two-axis control of a spin qubit, *Nat. Commun.* **15**, 1676 (2024).
- [39] T. Nakajima, Y. Kojima, Y. Uehara, A. Noiri, K. Takeda, T. Kobayashi, and S. Tarucha, Real-time feedback control of charge sensing for quantum dot qubits, *Phys. Rev. Appl.* **15**, L031003 (2021).
- [40] T. Nakajima, A. Noiri, K. Kawasaki, J. Yoneda, P. Stano, S. Amaha, T. Otsuka, K. Takeda, M. R. Delbecq, G. Allison, A. Ludwig, A. D. Wieck, D. Loss, and S. Tarucha, Coherence of a driven electron spin qubit actively decoupled from quasistatic noise, *Phys. Rev. X* **10**, 011060 (2020).
- [41] M. D. Shulman, S. P. Harvey, J. M. Nichol, S. D. Bartlett, A. C. Doherty, V. Umansky, and A. Yacoby, Suppressing qubit dephasing using real-time Hamiltonian estimation, *Nat. Commun.* **5**, 5156 (2014).
- [42] J. Park, H. Jang, H. Sohn, J. Yun, Y. Song, B. Kang, L. E. A. Stehouwer, D. D. Esposti, G. Scappucci, and D. Kim, Passive and active suppression of transduced noise in silicon spin qubits, *Nat. Commun.* **16**, 78 (2025).
- [43] G. Burkard, M. J. Gullans, X. Mi, and J. R. Petta, Superconductor-semiconductor hybrid-circuit quantum electrodynamics, *Nat. Rev. Phys.* **2**, 129 (2020).
- [44] A. Warren, E. Barnes, and S. E. Economou, Long-distance entangling gates between quantum dot spins mediated by a superconducting resonator, *Phys. Rev. B* **100**, 161303 (2019).
- [45] H. Ball, W. D. Oliver, and M. J. Biercuk, The role of master clock stability in quantum information processing, *npj Quantum Inf.* **2**, 16033 (2016).
- [46] X. Mi, J. V. Cady, D. M. Zajac, J. Stehlik, L. F. Edge, and J. R. Petta, Circuit quantum electrodynamics architecture for gate-defined quantum dots in silicon, *Appl. Phys. Lett.* **110**, 043502 (2017).
- [47] P. Harvey-Collard, G. Zheng, J. Dijkema, N. Samkharadze, A. Sammak, G. Scappucci, and Lieven M. K. Vandersypen, On-chip microwave filters for high-impedance resonators with gate-defined quantum dots, *Phys. Rev. Appl.* **14**, 034025 (2020).
- [48] Y. Tian, I. Grytsenko, A. Jennings, J. Wang, H. Ikegami, X. Zhou, S. Tamate, H. Terai, H. Kutsuma, D. Jin, M. Benito, and E. Kawakami, Data for “Nb-Ti-N nanowire resonators and spin-photon coupling prospects with electrons on solid neon,” <https://doi.org/10.5281/zenodo.18058729> (2025).
- [49] J. Gao, The physics of superconducting microwave resonators, Ph.D. thesis, California Institute of Technology, 2008.
- [50] E. Y. Andrei, *Two-Dimensional Electron Systems: On Helium and Other Cryogenic Substrates* (Springer, Netherlands, 1997).
- [51] Note that we use the convention $\exp(i\omega t)$ for the time dependence of fields. While the physics community commonly adopts $\exp(-i\omega t)$, the engineering convention $\exp(i\omega t)$ is used here for consistency with COMSOL.
- [52] D. M. Pozar, *Microwave Engineering* (John Wiley & Sons Inc., New York, 1998), 4th ed.
- [53] M. Gurioli, J. Martinez-Pastor, M. Colocci, C. Deparis, B. Chastaingt, and J. Massies, Thermal escape of carriers out of GaAs/Al_xGa_{1-x}As quantum-well structures, *Phys. Rev. B* **46**, 6922 (1992).
- [54] R. K. Swank and F. C. Brown, Lifetime of the excited F center, *Phys. Rev.* **130**, 34 (1963).
- [55] J. Gao, M. Daal, A. Vayonakis, S. Kumar, J. Zmuidzinas, B. Sadoulet, B. A. Mazin, P. K. Day, and H. G. Leduc, Experimental evidence for a surface distribution of two-level systems in superconducting lithographed microwave resonators, *Appl. Phys. Lett.* **92**, 152505 (2008).
- [56] M. Müller, T. Luschmann, A. Faltermeier, S. Weichselbaumer, L. Koch, G. B. Huber, H. W. Schumacher, N. Ubbelohde, D. Reifert, T. Scheller, *et al.*, Magnetic field robust high quality factor NbTiN superconducting microwave resonators, *Mater. Quantum Technol.* **2**, 015002 (2022).
- [57] A. Bruno, G. De Lange, S. Asaad, K. Van Der Enden, N. Langford, and L. DiCarlo, Reducing intrinsic loss in superconducting resonators by surface treatment and deep etching of silicon substrates, *Appl. Phys. Lett.* **106**, 182601 (2015).
- [58] R. Barends, N. Vercruyssen, A. Endo, P. De Visser, T. Zijlstra, T. Klapwijk, P. Diener, S. Yates, and J. Baselmans, Minimal resonator loss for circuit quantum electrodynamics, *Appl. Phys. Lett.* **97**, 023508 (2010).
- [59] J. R. Lane, D. Tan, N. R. Beysengulov, K. Nasyedkin, E. Brook, L. Zhang, T. Stefanski, H. Byeon, K. W. Murch, and J. Pollanen, Integrating superfluids with superconducting qubit systems, *Phys. Rev. A* **101**, 012336 (2020).