

Quasiballistic (In, Ga)As Josephson field-effect transistor with molybdenum electrodes

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Gate-tunable Josephson junctions with molybdenum superconducting contacts are reported. The Josephson field-effect transistors are based on an (In, Ga)As quantum well heterostructure with Mo contacts deposited on a layer of heavily doped (In, Ga)As by HCl cleaning and sputtering. We reproduced the measured data using a quasiballistic model that captures the gate dependence of the critical current and considers contributions from both Andreev states in the discrete spectrum and leaky states in the continuous spectrum. We report the observation of multiple Andreev reflections (MARs) up to the fourth order, and the measured temperature dependence of the MAR features is consistent with a model describing an effective induced superconducting gap beneath the contacts, approximately 55% of the gap in the Mo. The transistors are integrated into an rf-compatible geometry with Si₃N₄ spacers, enabling future implementation of highly scaled, compact, quantum-limited cryogenic amplifiers.

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I. INTRODUCTION

The integration of superconductors with semiconductors enables the fabrication of hybrid devices, combining the unique properties of superconductivity with the tunability of the semiconductor material. Growing interest in the potential applications of quantum technology has driven efforts to realize these hybrid superconductor-semiconductor devices, with focus on achieving high transparency of the superconductor/semiconductor interface [1]. Emerging applications of these devices include qubits for superconducting quantum computation [2–6], electrically tunable parametric amplifiers [7–9], single-photon detectors [10,11], and bolometers [12–14].

Narrow bandgap materials such as InAs [15] and InSb [16] have been studied for their high mobility, and (In, Ga)As hybrid devices have been investigated due to their potential for low-noise radio-frequency (rf) performance [17] and as a possible platform for topological quantum computing [1]. Significant efforts have been made to fabricate semiconductor systems coupled with aluminum, primarily due to its long bulk coherence length

ξ_0^{Al} of 1600 nm [18]. In the study reported in Ref. [17], *ex situ* Al contacts were deposited by e-beam evaporation on an (In, Ga)As two-dimensional electron gas (2DEG), resulting in poor interface transparency and high contact resistance on the order of $10^4 \Omega \mu\text{m}$.

One method of fabricating high-quality superconductor/semiconductor interfaces is by growing Al *in situ* by molecular beam epitaxy (MBE), which has been applied to InAs nanowires [19] and InAs/(In, Ga)As 2DEGs [20–22]. This procedure is believed to work due to the growth being *in situ*, preventing the formation of oxides, making a very clean, homogeneous interface. Moreover, the Fermi level at the InAs surface is pinned in the conduction band, also contributing to the resulting formation of transparent superconductor/semiconductor interfaces. Possible downsides of this method are related to cost, speed, complexity, and availability of MBE growth systems. It might be advantageous in some applications to find a simpler method, available to more researchers, that is faster and more cost-effective but still results in hybrid superconductor-semiconductor devices with high-quality interfaces.

Semiconductor-based Josephson weak links with a gate-tunable supercurrent have been explored as a circuit element in different circuit configurations [23] and for low-power Boolean logic applications [24]. More recent studies on 2DEG-based [8], nanowire-based [25], and graphene-based [7] Josephson parametric amplifiers, as an essential element for low-noise readout of quantum systems, have demonstrated gate-controllable resonance shifts of

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approximately 1 GHz, gain exceeding 20 dB, and noise performance near the quantum limit [7,25], providing concrete experimental benchmarks for the use of Josephson field-effect transistors (JoFET) as a source of electrically tunable and dissipation-free nonlinearity in rf amplification. These experimental demonstrations of gate-tunable Josephson parametric amplifiers suggest interesting alternatives for different amplification schemes.

Lattice-matched (In,Ga)As 2DEG systems grown on InP combine high electron mobility and low effective mass with the planar, uniform, and gate-tunable nature of the 2DEG, enabling fast, low-noise carrier transport and reproducible, scalable architectures for superconductor-semiconductor hybrid devices. We believe that by incorporation of the JoFETs into an rf-compatible geometry, such as a planar (In,Ga)As weak link with a T-shaped gate and low- k spacers, the gate can serve as a high-frequency pump port. This enables on-chip integration of electrical pumping via the gate, reducing the need for large external pump signals and allowing fast modulation of the critical current, providing a route toward compact quantum-limited amplifiers.

In this paper, we investigate the superconducting properties of Mo-(In,Ga)As-Mo Josephson field-effect transistors with sputtered Mo contacts on top of n^+ (In,Ga)As, achieving a relatively low contact resistance of 30 Ω μ m, with the *ex situ* fabrication process being promising for (In,Ga)As devices. This is, to our knowledge, the first time molybdenum has been used to fabricate Josephson junctions. We observe an induced superconducting gap of approximately 55% of the Mo gap and a high value of $I_{\text{ex}}R_N/\Delta \approx 1.5$, indicating good quality of the Mo/ n^+ (In,Ga)As interface. We also show that the gate dependence of the critical current can be accurately reproduced with a quasiballistic model that considers contributions from both discrete Andreev states and leaky states in the continuous spectrum. The model uses only the mobility from the normal-state resistance and the induced gap from multiple Andreev reflection (MAR) as input parameters.

II. METHOD

A. Device fabrication

The JoFETs were fabricated on a wafer consisting of epitaxial layers of III-V alloys on top of a semi-insulating InP substrate, in a layer stack according to Table I. The 2DEG consists of layers 4, 5, and 6, with maximum electron carrier concentration in layer 5.

The first step of the sample fabrication was removal of the native oxide in HCl-H₂O (1:20), followed by dc magnetron sputtering of 30 nm Mo at 0.3 mTorr on top of a layer of heavily doped (In,Ga)As. The use of an n^+ (In,Ga)As cap together with Mo is a proven technique for forming Ohmic contacts with low contact resistance

TABLE I. Sample layer stack. SI, semi-insulating.

Layer	Material	x	Thickness ($\text{\AA}$)	Si (cm^{-3})	Function
10	In _(x) Ga _{1-x} As	0.63	50	4.0×10^{19}	Drain
9	In _(x) Ga _{1-x} As	0.53	200	4.0×10^{19}	Drain
8	InP		60	1.0×10^{19}	Etch stop
7	InP		10		Etch stop
6	In _(x) Ga _{1-x} As	0.53	30		Channel
5	In _(x) Ga _{1-x} As	0.80	50		Channel
4	In _(x) Ga _{1-x} As	0.53	40		Channel
3	In _(x) Al _{1-x} As	0.52	30		Spacer
2	In _(x) Al _{1-x} As	0.52	30	1.0×10^{19}	δ doping
1	In _(x) Al _{1-x} As	0.52	1900		Buffer
0	SI InP				Substrate

[26]. Higher doping concentrations increase the Fermi level and increase the number of conducting modes. An increased carrier density allows easier carrier injection at the metal/semiconductor interface, leading to a reduction in contact resistance.

After sputtering, deposition of a 30-nm-thick Si₃N₄ spacer using plasma-enhanced chemical vapor deposition followed. Spacers of Si₃N₄ are used to improve future rf performance of the devices by reducing the parasitic overlap capacitance of the parallel-plate-like capacitors between the source and drain contact layers and the gate metal. The use of spacers also contributes to reducing gate leakage, which is crucial when thin oxides necessary for achieving high transconductance are used.

Openings in the Mo film were then defined by electron-beam lithography (EBL) and F-based inductively coupled plasma reactive ion etching (ICPRIE) of the Si₃N₄ and Mo layers. The n^+ (In,Ga)As between the Mo electrodes was then removed in a wet etch using H₃PO₄-H₂O₂-H₂O (1:1:25), and InP layers 7 and 8 were removed in a wet etch using HCl-H₂O (1:1). The wet etches are isotropic, giving rise to an undercut and slightly greater gate lengths L_g than the openings in the Mo layer. The devices were then electrically isolated by a mesa with use of ultraviolet lithography (UVL), ICPRIE, and wet etches into layer 1. Next, a thin gate oxide of approximately 10 nm was deposited by atomic layer deposition (ALD), including *in situ* surface cleaning by prepulsing of trimethylaluminum (TMA) (five cycles). This reduces the presence of Ga—O and As—O bonds [27] and restores surface As dangling bonds to a bulklike bonding configuration, thereby reducing the density of interface traps [28]. The prepulsing of TMA was followed by deposition of an Al₂O₃/HfO₂ bilayer gate oxide (7/100 cycles) at 300/120 $^{\circ}$ C, which results in good oxide/semiconductor interface quality and a low equivalent oxide thickness. We use a thinner gate oxide than in comparable studies [21,22,29], which enhances electrostatic control over the supercurrent. This design is essential for achieving very short gate length scaling,

akin to that in high-performance field-effect transistors [30]. Gate length scaling in conjunction with high gate coupling can lead to both strong modulation of the supercurrent with gate voltage and at the same time a small gate capacitance, which can lead to active superconducting microwave devices. The next step in the fabrication process was to define 30-nm-thick TiN gate electrodes by ALD (300 cycles) at 250 °C, EBL, and F-based ICPRIE. After this, the high- k gate oxide and the Si₃N₄ spacer outside the gate region were removed by buffered oxide etch. Finally, a UVL lift-off process and electron-beam evaporation of Ti(5 nm)/Pd(5 nm)/Au(200 nm) was used to define the measurements pads. Figure 1(a) shows a schematic and Fig. 1(b) shows a false-color scanning electron micrograph of the final device.

B. Model of induced superconducting gap

The charge transport through a superconductor/semiconductor interface is mediated by the mechanism of Andreev reflection, which occurs at the two boundaries indicated by the dashed blue and green lines in Fig. 1(a). The superconductivity in this planar geometry can be modeled according to the extended version of the Octavio-Blonder-Tinkham-Klapwijk (OBTK) model denoted $SNcNS$ [31], where N is the normal part of the 2DEG underneath the superconducting Mo and c is the constricted part of the 2DEG of length L_g without Mo on top. Because of MAR, the N part can be considered to gain an induced superconducting gap Δ^* , smaller than the gap of the Mo Δ_{Mo} . The magnitude of the induced gap is set by the effective transparency of the Mo/2DEG interface, indicated by the horizontal dashed blue lines in Fig. 1(a). The induced gap then gives rise to MAR at the effective 2DEG/uncovered (In,Ga)As interface, indicated by the vertical dashed green lines. The variation of the

superconducting gap $\Delta(z)$ in the growth direction z is indicated in the right schematic in Fig. 1(a). Δ^* extends for a length $d_N \ll \xi_N$ in the 2DEG, where ξ_N is the normal-state coherence length in the semiconducting region covered by Mo.

In this paper, two JoFETs with gate lengths $L_g = 140$ nm and $L_g = 320$ nm and device width $W = 4.9$ μm are studied. The devices were characterized by current-driven four-terminal measurements at cryogenic temperatures in a dilution refrigerator with base chamber temperature $T \approx 15$ mK.

III. RESULTS AND DISCUSSION

A. Normal resistance

Figure 2(a) shows the voltage drop V over the junction as a function of current bias I for the device with $L_g = 320$ nm and gate voltage $V_{gs} = 0.25$ V, indicating the critical current I_c and the excess current I_{ex} . The critical current is the maximum supercurrent supported in the junction, and for higher currents the system switches to the resistive state. Experimentally, there is a difference between the critical current and the observed switching current, which is lower, since temperature and noise can cause the junction to switch earlier [21]. In this paper, for simplicity we assume that the two currents are equal, as they would be for an ideal device, which thereby can lead to a slight underestimation of I_c . The Andreev reflection process that carries the supercurrent across the junction gives rise to the excess current $I_{ex} = I - V/R_N$, which can be extracted as the intercept of $V = 0$ with a linear fit from the high-current normal regime, taken at $V \gg \Delta^*/q$ [32], where q is the elemental charge, as indicated in Fig. 2(a). The slope of the linear fit at large V corresponds to R_N , which is the normal resistance of the junction in the absence of superconductivity. The magnitude of the induced gap Δ^* was found to

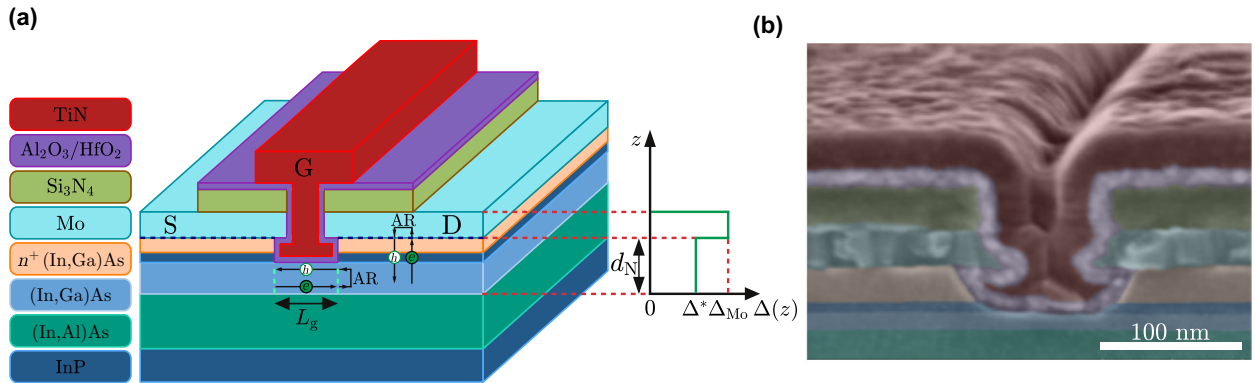


FIG. 1. (a) The JoFET (not to scale) with source (S), drain (D), gate (G), and gate length L_g . Andreev reflection (AR) can happen at both the horizontal (blue) and the vertical (green) interfaces indicated by dashed lines. The schematic to the right shows the variation of the superconducting gap $\Delta(z)$ in the growth direction z , where Δ^* is the induced gap, which extends a distance d_N underneath the Mo contacts, and Δ_{Mo} denotes the superconducting gap of the Mo. (b) False-color scanning electron micrograph of the device.

be 110 μV , and is discussed in Sec. III C concerning the experimental measurements of MAR.

Figure 2(b) shows the reduced normal resistance $R_N^* = R_N - 2R_c$ times the width W versus V_{gs} . Within the framework of the $SNcNS$ model, R_N^* is particularly suitable because it separates the contributions of the channel and the interface. The normal-state resistance R_N^* reflects the transmission through the constriction and provides a measure of the transport efficiency in the 2DEG, whereas the contact resistance R_c is associated with the superconductor/semiconductor interface transparency and affects the magnitude of the induced superconducting gap Δ^* . Therefore, consideration of R_N^* instead of the standard figure of merit R_N , provides a more accurate benchmark for understanding the behavior of superconductor/2DEG/superconductor junctions.

The contact resistances, $2R_c^{140\text{ nm}} = 140\ \Omega\ \mu\text{m}$ and $2R_c^{320\text{ nm}} = 30\ \Omega\ \mu\text{m}$, were determined from $2R_c = (R_N - R_N^{\text{sim}})|_{V_{gs}=1\text{ V}}$, with R_N^{sim} calculated from a quasiballistic model considering the gate-dependent transport transmission probability per conducting mode [33]

$$\tau = \frac{\lambda}{\lambda + L_g}, \quad (1)$$

where λ is the average mean free path related to the long-channel mobility μ defined as [33]

$$\lambda = n_s q \mu \frac{h}{2q^2} \frac{W}{M}, \quad (2)$$

where n_s is the electron carrier concentration, h is Planck's constant, and M is the number of conducting modes in the channel. The transport in the junctions appears quasiballistic. A proportional increase in channel resistance with gate length would be expected for a diffusive junction; however, the observed increase is only by a factor of 1.8 (from 210 to 380 $\Omega\ \mu\text{m}$), while the gate length increases by a factor of 2.3 (from 140 to 320 nm), indicating quasiballistic transport. For a fully ballistic junction there would be no expected increase in channel resistance with increasing gate length. While this is an indication of quasiballistic behavior, more statistics are needed to strongly confirm the assumption. The measured R_N^* in Fig. 2(b) shows good agreement with simulated data. Details of the simulation can be found in Appendix A. To confirm the R_N values from the high- V measurements, measurements of R_N were also performed at a large perpendicular magnetic field that exceeds the critical magnetic field ($B_c = 90\text{ mT}$) of the Mo contacts. This method results in the same values of R_N , which can be seen in the inset in Fig. 2(b).

Figure 2(c) presents the gate-dependent mobility, which follows the expected behavior of a 2DEG at low temperature [34,35]. The mobility increases with gate voltage at low carrier concentrations, reaches a maximum, and then

decreases at higher carrier concentrations. At low carrier concentrations, the mobility is limited primarily by scattering from remote impurities. As the carrier concentration increases, enhanced screening suppresses this scattering mechanism, resulting in higher mobility. At high carrier concentrations, the mobility decreases. The charge carriers move closer to the high- k oxide/III-V interface, leading to increased scattering rates, primarily due to scattering from charged surface defects and surface roughness in the near-surface quantum well. Detailed analyses of the carrier density-dependent mobility and scattering mechanisms were performed in previous studies on devices fabricated from the same wafer stack [34,35]. The mobility obtained agrees well for both devices, with a peak mobility of 3200 $\text{cm}^2/\text{V s}$. With a maximum mean free path $\lambda = 110\text{ nm} < L_g$, this also indicates that the devices are operating in the quasiballistic regime. The low peak mobility originates from the quantum well being placed directly in contact with the high- k oxide. Buried quantum wells can strongly improve the mobility, at the cost of a weaker gate coupling.

In the 2DEG devices, several conducting subbands indexed by u are present, each characterized by different Fermi velocities v_F and coherence lengths $\xi_N = \hbar v_F / 2\Delta^*$ as shown in Fig. 2(d) for the device with $L_g = 140\text{ nm}$. Most of the subbands have coherence lengths exceeding the gate length, placing them in the short-junction regime. In contrast, the subbands with the slowest Fermi velocities have coherence lengths shorter than the gate length, and are thus better described within the long-junction regime.

B. Critical current

Figure 3(a) shows the critical current. For negative gate biases, the critical current approaches zero, and can be well controlled by a small applied gate voltage. We report a maximum value of the critical current of 2.4 μA for the shorter device and 820 nA for the longer device. Figure 3(b) shows $I_c R_N^* / \Delta^*$ as a function of gate voltage. The maximum values are 1.3 ($I_c R_N / \Delta = 0.95$) and 0.61 ($I_c R_N / \Delta = 0.36$) for the two devices. These values can be compared with the maximum theoretical value of π for a fully transparent short ($\xi_N \gg L$) ballistic junction when $\lambda \gg L$, and the maximum theoretical value of $1.32\pi/2 \approx 2.07$ for a short diffusive junction when $L \gg \lambda$, where L is the effective length of the junction [36]. The $I_c R_N^* / \Delta^*$ values for the devices are 41% and 19% of the ballistic limit, respectively, and 63% and 29% of the diffusive limit, respectively, with the higher values for the shorter device indicating its more ballistic transport characteristics. Compared with other reported values, the $I_c R_N^* / \Delta^*$ values obtained here are lower than earlier results of approximately 70% in Refs. [21,22], with the lower value originating from a smaller λ/L_g ratio and lower

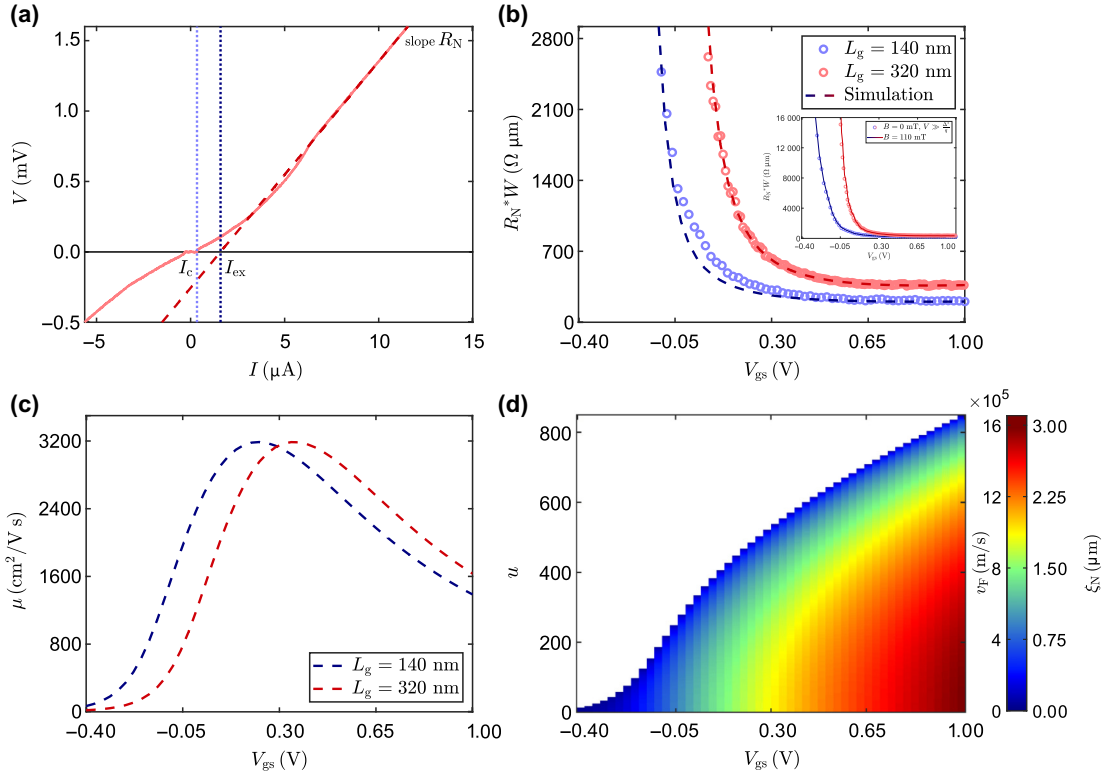


FIG. 2. (a) Voltage V across the junction as a function of current bias I for the device with $L_g = 320$ nm with applied gate bias $V_{gs} = 0.25$ V. The dashed red line with a slope of R_N shows a linear fit taken at $V \gg \Delta^*/q$. The intercept of the linear fit with $V = 0$ gives the excess current I_{ex} , indicated by a dashed dark blue line. The critical current I_c is the maximum supercurrent and is indicated by a dashed light blue line. (b) Experimental and simulated data for the reduced normal resistance $R_N^* = R_N - 2R_c$, where R_c is the contact resistance, times the width W as a function of gate voltage V_{gs} . The inset in (b) shows R_N^*W as a function of gate voltage, extracted in measurements with magnetic field $B = 0$ mT and $B = 110$ mT, respectively. (c) Gate-dependent mobility μ used to simulate R_N^* as a function of gate voltage. (d) Fermi velocity v_F and coherence length ξ_N as a function of gate voltage and subband index u for the device with $L_g = 140$ nm.

mobility. For a fully ballistic device, the critical current would theoretically approach $20 \mu\text{A}$.

To try to understand the dependence of the critical current on the gate voltage in the quasiballistic devices within the $SNcNS$ model, we modeled the critical current using a basic quasiballistic one-dimensional (1D) model [37] valid for both long and short junctions, reflecting the presence of both short and long subbands in the devices. For a specific gate voltage, the critical current is obtained by summing over all subband contributions, taking into account the different Fermi velocity of each subband as shown in Fig. 2(d). Instead of a ballistic constriction, we used a straightforward model where a point impurity potential $V_a(x) = V_s \delta(x - a)$ is placed in the center of the constriction at position $a = L_g/2$, where V_s is the impurity amplitude, set to reproduce the measured value of R_N^* , thus simulating a quasiballistic device. To evaluate the sensitivity of the critical current to the impurity position, additional simulations were performed with the impurity placed at the beginning of the constriction with $a = 0$. For instance, at $V_{gs} = 0.2$ V and $L_g = 140$ nm, the maximum percentage deviation relative to the configuration

with $a = L_g/2$ was only 0.08%, with absolute differences on the order of nanoamperes. This indicates that, despite the simplicity of the model, the critical current is negligibly affected by the exact impurity position within the constriction.

The model takes into account both the current contribution I_{discrete} carried by the discrete Andreev levels for energies E inside the superconducting energy gap, $-\Delta^* < E < \Delta^*$, and the contribution $I_{\text{continuous}}$ from scattering states for energies outside the gap, $E > \Delta^*$ and $E < -\Delta^*$. A schematic of the energy landscape is shown in Fig. 3(c), where the induced superconducting gap is modeled as a step function

$$\Delta^*(x) = \begin{cases} \Delta^* e^{i\phi_1}, & x < 0, \\ 0, & 0 < x < L_g, \\ \Delta^* e^{i\phi_2}, & x > L_g, \end{cases} \quad (3)$$

where ϕ_1 and ϕ_2 are the superconducting phases associated with each induced superconductor on either side of the constriction and x denotes the channel direction. Details of the simulation can be found in Appendix B.

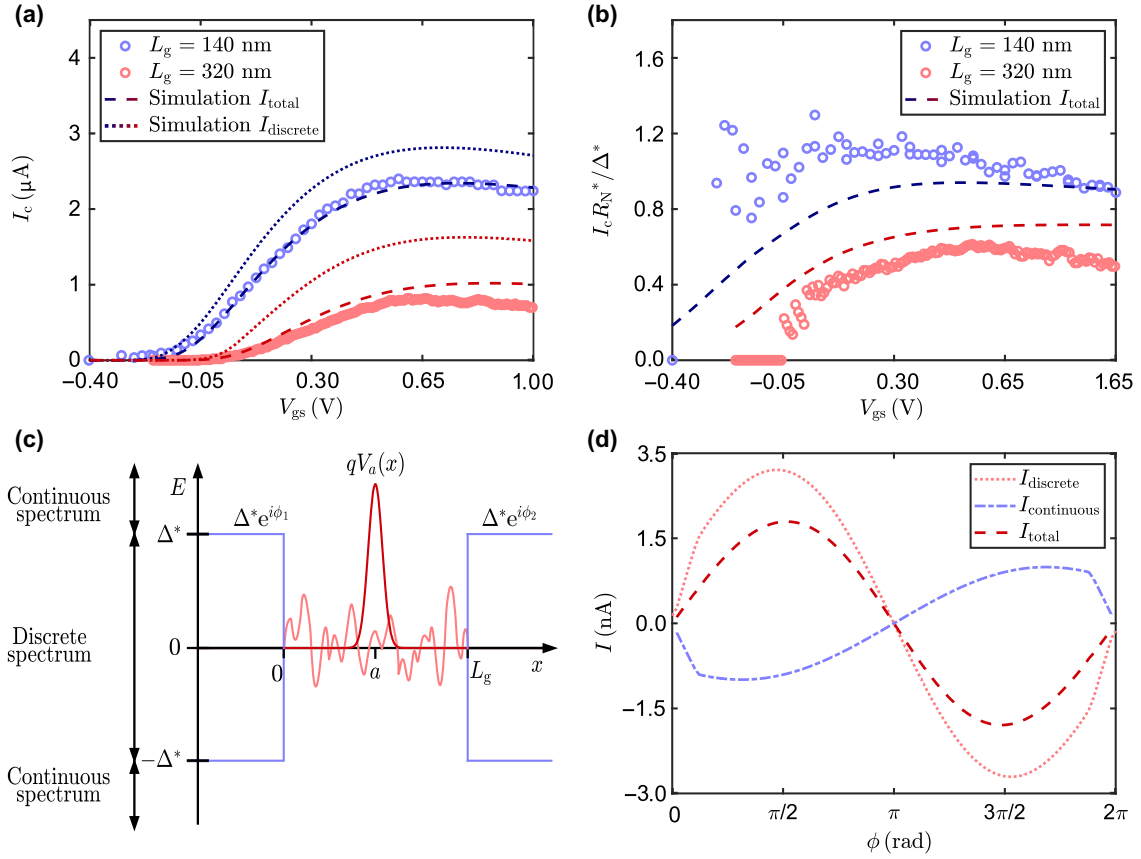


FIG. 3. (a) Experimental and simulated data for the critical current I_c as a function of gate voltage V_{gs} . The simulated data show the critical current when only the discrete contribution is included and when both contributions are included. The maximum experimental values of the critical current are $2.4 \mu\text{A}$ for the shorter device and 820 nA for the longer device. (b) Experimental and simulated data for $I_c R_N^* / \Delta^*$ as a function of gate voltage. The maximum experimental values of $I_c R_N^* / \Delta^*$ are 1.3 for the shorter device and 0.61 for the longer device. (c) Energy landscape of the JoFET showing the energy E as a function of position x . The model considers an induced superconducting gap Δ^* modeled as a step function (purple lines) with a point impurity potential $V_a(x)$ in the center of the constriction (red lines). More realistically, energy fluctuations in the constriction are indicated by pink lines. (d) Current contributions I_{discrete} and $I_{\text{continuous}}$ and the total current I_{total} as a function of the superconducting phase difference ϕ .

Figure 3(d) shows I_{discrete} , $I_{\text{continuous}}$, and the total current I_{total} as a function of the superconducting phase difference $\phi = \phi_2 - \phi_1$ for a subband of the device with $L_g = 140 \text{ nm}$. For the devices studied here, I_{discrete} and $I_{\text{continuous}}$ are out of phase with each other, and consideration of the continuous spectrum is needed to obtain a more correct amplitude of the supercurrent.

Interestingly, the simulation in Fig. 3(a) shows excellent agreement for the device with $L_g = 140 \text{ nm}$ despite the simplicity of the multisubband quasiballistic 1D model. For the longer $L_g = 320 \text{ nm}$ device, the agreement is reasonable, somewhat overestimating the current, possibly indicating a transition toward a more diffusive transport model [38]. It should be noted that no free fitting parameters have been used, with the transmission obtained from the measurement of R_N^* , and the induced gap Δ^* obtained from MAR measurements. The critical current begins to saturate at V_{gs} around 0.3 V , approximately coinciding with the mobility in Fig. 2(c)

reaching its maximum value. At higher gate voltages, the mobility decreases, while the carrier concentration continues to increase linearly [$n_s \approx C_g(V_{gs} - V_T)$], leading to reduced transmission per mode but an increased number of conducting modes. This combination is likely the main factor underlying the observed saturation or slight decrease of the critical current at higher gate voltage overdrives.

Within the used model, the simulated data when only the discrete spectrum is taken into account overestimate the critical current for both junctions, and the difference is more pronounced for the longer device. This indicates that inclusion of the leaky Andreev states gives a more accurate representation of the critical current and is something that needs to be considered for accurate modeling if the gate length is approaching the coherence length. In the multisubband 1D model, for a certain gate voltage overdrive, the current is carried by both long and short subbands, as indicated in Fig. 2(d).

The measured and simulated values of $I_c R_N^*/\Delta^*$ shown in Fig. 3(b) again show good agreement for both gate lengths. Deviations in the $I_c R_N^*$ product are largely due to the differences between the measured and the simulated R_N^* obtained from the mobility fit [see Figs. 2(b) and 2(c)]. The simulated I_c and R_N^* both rely on this fit, but the deviations in R_N^* are more apparent, and therefore dominate the deviations observed in $I_c R_N^*$. In the model, for example, both devices are assumed to have the same subthreshold swing and mobility, apart from a shift in V_T . Performing optimized fits separately for each device would result in more accurate simulations of the $I_c R_N^*$ products. From the modeling, we can conclude that the reduced $I_c R_N^*/\Delta^*$ comes in part from the quasiballistic nature of the devices and in part from the effect of the continuum states. While the simple 1D $SNcNS$ model agrees well with the measured data for the short device, one should note that the model assumes fully phase coherent transport, no intersubband scattering,

and an average transmission equal for all subbands. For a two-dimensional fully ballistic junction with planar contacts, the more advanced model in Ref. [39] can be used to better capture the physics of the device, especially the temperature dependence.

C. Multiple Andreev reflection and excess current

The Andreev reflection mechanism is responsible for the flow of supercurrent through a superconductor/normal conductor/superconductor junction. In the voltage state, we observe multiple Andreev reflections, which manifest themselves as a series of peaks in the differential conductance at fixed voltage positions, and can be seen in Fig. 4(a) for the device with $L_g = 140$ nm. The MAR features are indicated by dashed red lines, and appear to be almost independent of the gate voltage. For very low gate voltages V_{gs} below -0.3 V, the channel is depleted, with almost no current flow and very high resistance.

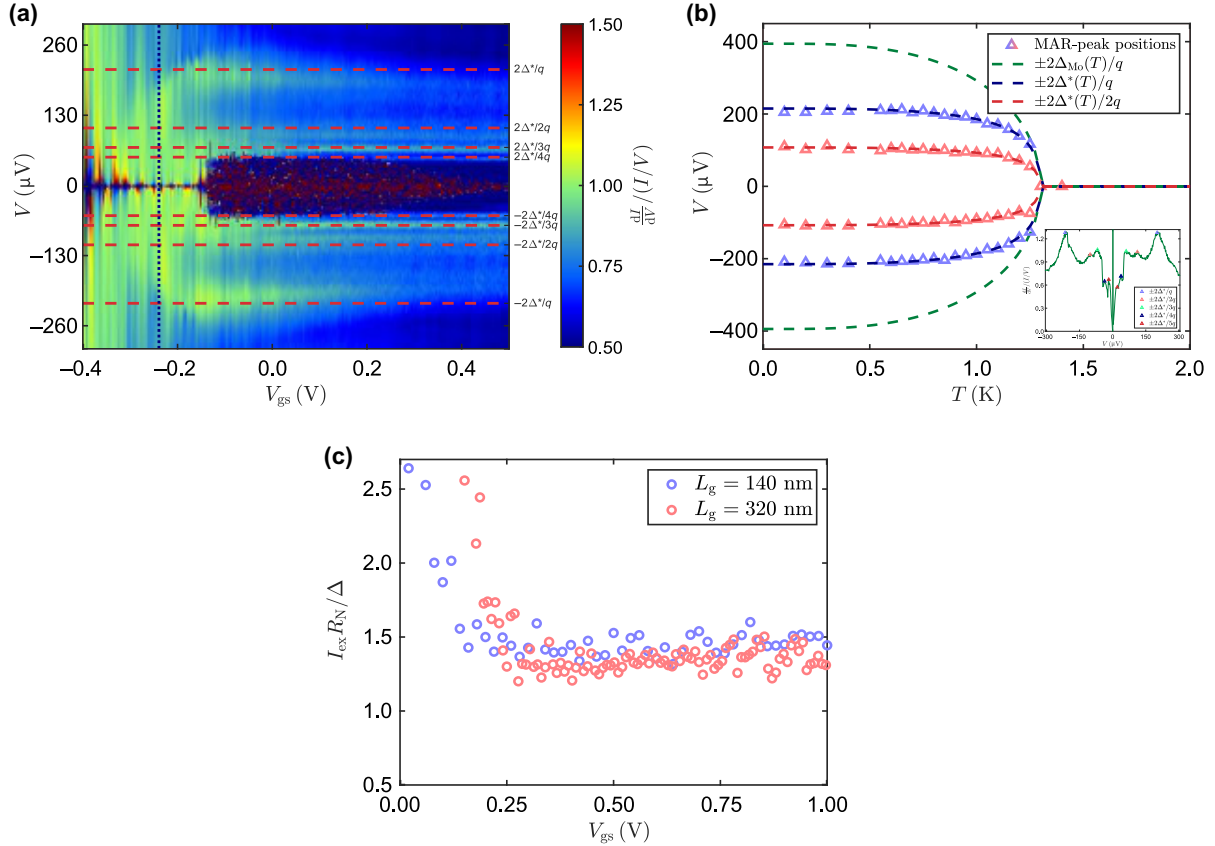


FIG. 4. (a) Normalized differential conductance $(dI/dV)/(I/V)$ as a function of junction voltage V and gate voltage V_{gs} for the device with $L_g = 140$ nm. MAR peaks up to fourth order are indicated by dashed red lines and occur at multiples of an induced gap $\Delta^* = 110$ μ V and appear almost independent of V_{gs} . (b) MAR-peak positions for $n = \pm 1, \pm 2$ as a function of temperature T for $V_{gs} = -0.24$ V, corresponding to the dashed blue line in (a). Predictions of the temperature behavior of the induced gap and the Mo gap are also shown. The inset shows the normalized differential conductance as a function of voltage for $T = 300$ mK. (c) $I_{ex} R_N / \Delta$ as a function of gate voltage. $I_{ex} R_N / \Delta$ appears rather constant at around 1.5 for the device with $L_g = 140$ nm and at around 1.4 for the device with $L_g = 320$ nm.

The increased channel resistance introduces a significant voltage drop outside the junction region, making the voltage-bias measurements unreliable.

According to the predictions from the OBTK theory [40,41], observation of peaks in the differential conductance is expected to occur at submultiples of the gap voltage $V = 2\Delta^*/qn$, $n = \pm 1, \pm 2, \pm 3, \dots$. The position of the MAR features is not expected to change substantially with gate voltage [42,43]. We measure MAR peaks up to order $n = 4$. The peaks occur at multiples of $\Delta^* = 110 \mu\text{V} \approx 55\%$ of $\Delta_{\text{Mo}} = 1.764k_{\text{B}}T_{\text{c}} \approx 200 \mu\text{V}$, where k_{B} is the Boltzmann constant and $T_{\text{c}} = 1.3 \text{ K}$ is the critical temperature obtained from four-terminal measurements on planar Mo films. The induced gap of approximately 55% of the Mo gap can be compared with induced gaps observed in other studies; for example, an induced gap in InAs/(In, Ga)As 2DEG Josephson junctions of approximately 76% [22], approximately 81% with *in situ* epitaxial Al contacts [20], and approximately 9.1% with Nb contacts [1], and approximately 2.4% [44] in GaAs 2DEGs with NbN contacts.

The MAR-peak positions of order $n = \pm 1, \pm 2$ are shown in Fig. 4(b) as a function of temperature T for $V_{\text{gs}} = -0.24 \text{ V}$ corresponding to the dashed blue line in Fig. 4(a). The temperature dependence of MAR peaks for $V_{\text{gs}} = -0.12 \text{ V}$ can be found in Appendix C. The inset in Fig. 4(b) shows the voltage positions of the MAR peaks for $T = 300 \text{ mK}$. The MAR-peak positions agree very well with the model of an induced superconducting gap, whose temperature dependence self-consistently is given by [31]

$$\Delta^*(T) = \frac{\Delta_{\text{Mo}}(T)}{1 + \gamma_{\text{B}} \sqrt{\Delta_{\text{Mo}}^2(T) - \Delta^{*2}(T)/\pi k_{\text{B}}T_{\text{c}}}} \quad (4)$$

under the assumption that $d_{\text{N}} \ll \xi_{\text{N}}$. $\Delta_{\text{Mo}}(T)$ was determined from BCS theory and γ_{B} is a dimensionless parameter that determines the difference between the induced gap and the Mo gap and is a measure of the boundary transparency of the superconductor/semiconductor interface, indicated by the horizontal dashed blue lines in Fig. 1(a). $\gamma_{\text{B}} = 0$ indicates a perfectly transparent interface. We find $\gamma_{\text{B}} = 1.8$, potentially limited by the presence of the InP etch stop layer and the highly disordered n^+ (In, Ga)As cap, which agrees reasonably well with estimates of $\gamma_{\text{B}} \approx 2\text{--}3$ from the R_{c} values obtained following the method described in Ref. [31]. Finding an optimized thickness of the n^+ layer and removing the InP layer may result in a lower γ_{B} value and therefore a larger induced gap. The measurements of the MAR peaks clarify that there is a difference in the temperature dependence between the induced gap and the conventional BCS form. Note that for both gate biases investigated, the MAR peaks are well fitted with Eq. (4) with the same values of γ_{B} and Δ^* .

When one is considering the interface quality, $I_{\text{ex}}R_{\text{N}}/\Delta$ is an important figure of merit, since a high interface transparency corresponds to a high probability of Andreev reflection and therefore a high excess current [21]. $I_{\text{ex}}R_{\text{N}}/\Delta$ is related to the dimensionless parameter Z , which is a measure of the superconductor/semiconductor interface transparency. For a fully transparent interface, $Z = 0$ and corresponds to $I_{\text{ex}}R_{\text{N}}/\Delta = 8/3 \approx 2.6$ for a ballistic junction [45] and $I_{\text{ex}}R_{\text{N}}/\Delta = (\pi^2/4) - 1 \approx 1.467$ for a diffusive junction [46]. Figure 4(c) shows $I_{\text{ex}}R_{\text{N}}/\Delta$ versus gate voltage for both devices. $I_{\text{ex}}R_{\text{N}}/\Delta$ seems rather independent of gate voltage, and stays at around 1.5 for the device with $L_{\text{g}} = 140 \text{ nm}$ and at slightly lower value for the longer device. The extraction of I_{ex} becomes sensitive to noise at lower voltages when the resistance in the channel is high. $I_{\text{ex}}R_{\text{N}}/\Delta \approx 1.5$ corresponds to $Z \approx 0.44$, around 56% of the ballistic limit and 102% of the diffusive limit, similar to the values reported in Ref. [21] with *in situ* epitaxial Al contacts. A high value of $I_{\text{ex}}R_{\text{N}}/\Delta$ indicates good transparency of the Mo/2DEG interface. Both a decrease [47] and no significant change [21] in $I_{\text{ex}}R_{\text{N}}$ with junction length were previously reported. An $I_{\text{ex}}R_{\text{N}}$ product independent of gate length is an indication of a highly transparent interface [21]. $I_{\text{ex}}R_{\text{N}}$ in Fig. 4(c) seems rather independent of gate voltage, at least for higher voltages, which was observed in previous studies [17,21]. This emphasizes that the $I_{\text{ex}}R_{\text{N}}$ product depends primarily on the interface transparency and is not affected by the change in carrier concentration in the 2DEG [21].

D. Joule heating

The normalized differential conductance $(dI/dV)/(I/V)$ for the device with $L_{\text{g}} = 140 \text{ nm}$ is shown in Fig. 5(a). The device is measured in a current bias configuration, and since the channel resistance varies with the gate voltage, V will be very large for low V_{gs} and very small for high V_{gs} . Two sharp dips in the differential conductance can be seen symmetrically around $V = 0 \text{ V}$, indicated by the red lines in Fig. 5(a), which corresponds to a power P of 4 nW. These dips disappear if the superconductivity is quenched by application of a perpendicular magnetic field B of $110 \text{ mT} > B_{\text{c}}$, as shown in Fig. 5(b). The dips can be explained by a localized breakdown of the superconductivity in the vicinity of the semiconducting region. This phenomenon occurs due to Joule heating [48] generated by resistive dissipation when current flows through regions with finite resistance. The thermal effects give rise to an increased localized temperature near the semiconductor. When the temperature exceeds the critical temperature, this leads to a localized suppression of the superconductivity. The transition from the superconducting state to the resistive state causes an abrupt change in the current-carrying mechanism, leading to a noticeable dip in the differential conductance [48].

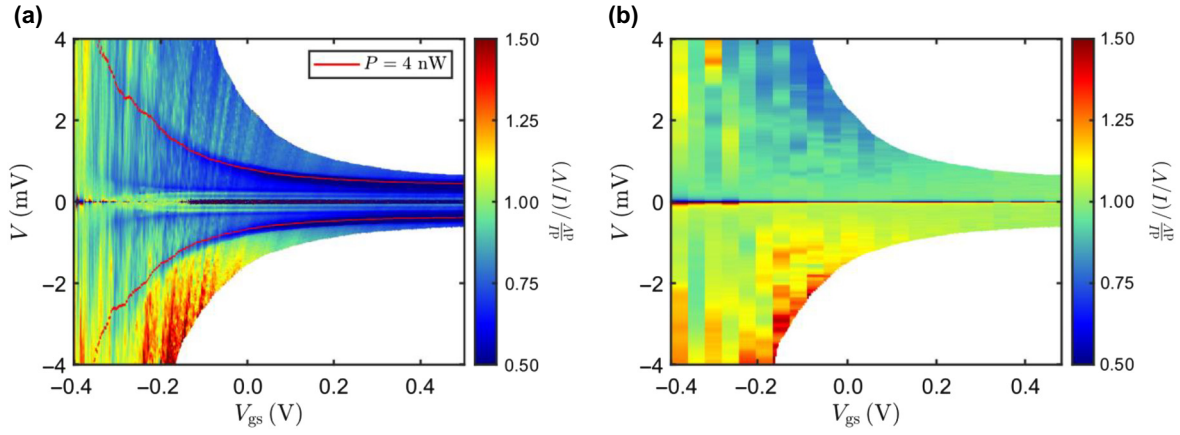


FIG. 5. (a) Normalized differential conductance $(dI/dV)/(I/V)$ as a function of measured junction voltage V and gate voltage V_{gs} . The red line indicates two dips in the normalized differential conductance due to localized suppression of the superconductivity caused by Joule heating. The dips occur at power P of approximately 4 nW. (b) Normalized differential conductance as a function of measured junction voltage and gate voltage with applied magnetic field $B = 110$ mT to quench the superconductivity. No dips can be seen.

IV. CONCLUSIONS

In conclusion, we have fabricated gate-tunable Josephson junctions using molybdenum superconducting contacts. The Mo-(In, Ga)As-Mo devices are integrated into an rf-compatible geometry, which serves as a promising candidate for a future platform for quantum-limited cryogenic amplifiers. We show good agreement between experimental and simulated critical currents using a straightforward quasiballistic model. We also report an induced superconducting gap $\Delta^* = 110 \mu\text{V}$ approximately 55% of the size of the Mo gap and high values of $I_{ex}R_N/\Delta \approx 1.5$. This indicates that it is possible to achieve a good-quality Mo/ n^+ (In, Ga)As interface by a simple method of HCl cleaning and sputtering.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [49].

APPENDIX A: QUASIBALLISTIC MODEL

We have used a quasiballistic model where the simulated normal resistance can be expressed as

$$R_N^{\text{sim}} = \frac{1}{M} \frac{h}{2q^2} \frac{1}{\tau}, \quad (\text{A1})$$

where τ is the transport transmission probability per conducting mode, which is related to the gate-dependent

mean-free path λ as [33]

$$\tau = \frac{\lambda}{\lambda + L_g}, \quad (\text{A2})$$

where L_g is the gate length. The mean free path is related to the gate-dependent mobility μ according to [33]

$$\lambda(n_s) = n_s q \mu(n_s) \frac{h}{2q^2} \frac{W}{M}, \quad (\text{A3})$$

$$\mu(n_s) = 0.0165 \left(\frac{1}{(20n_s/10^{16})^{-0.6}} + \frac{1}{(0.006n_s/10^{16})^{-1.8}} \right), \quad (\text{A4})$$

where W is the device width, n_s is the electron carrier concentration, h is Planck's constant, q is the elemental charge, and M is the number of conducting modes in the channel, which can be calculated according to

$$M = \frac{W}{(\pi/k_F)} = W \sqrt{\frac{2n_s}{\pi}}, \quad (\text{A5})$$

where k_F is the magnitude of the Fermi wave vector. An approximate expression for the carrier concentration valid for both $V_{gs} \gg V_T$ and $V_{gs} \approx V_T$ is [50]

$$n_s = \frac{C_g m \phi_t}{q} \ln \left(1 + e^{\frac{V_{gs} - V_T}{m \phi_t}} \right), \quad (\text{A6})$$

where V_{gs} is the gate voltage, $V_T (V_T^{140 \text{ nm}} = -0.15 \text{ V}, V_T^{20 \text{ nm}} = -0.02 \text{ V})$ is the threshold voltage, $\phi_t = k_B T/q$ is the thermal voltage, where k_B is the Boltzmann constant and T is the temperature, and m is the subthreshold

coefficient according to

$$m = \frac{S}{\phi_1 \ln 10}, \quad (\text{A7})$$

where $S = 48 \text{ mV/dec}$ is the subthreshold swing. In Eq. (A6), $C_g = 0.0065 \text{ F/m}^2$ is the total intrinsic gate capacitance, calculated with the oxide capacitance $\epsilon_{\text{ox}} = 16.4$ obtained from reference C - V measurements.

APPENDIX B: CRITICAL CURRENT MODEL

We have simulated the critical current by following the work of Bagwell [37]. The model considers both the current contribution I_{discrete} , carried by the discrete Andreev levels for energies E inside the superconducting energy gap Δ^* , $-\Delta^* < E < \Delta^*$, and the contribution $I_{\text{continuous}}$ from the continuous spectrum for energies outside the gap, $E > \Delta^*$ and $E < -\Delta^*$. The total critical current can then be calculated according to

$$I_c = \max \left(\sum_{u,n} (I_{\text{discrete}}(u, \phi) + I_{\text{continuous}}(u, \phi)) \right), \quad (\text{B1})$$

where $n = 0, \pm 1, \pm 2, \dots$ represents the index of the Andreev level, u is the index of the conducting mode, and $\phi = \phi_2 - \phi_1$ is the superconducting phase difference.

The current contribution from the discrete energy spectrum can be determined by identifying the Andreev levels $E_n^\pm(u, \phi)$ satisfying [37]

$$2\cos^{-1} \left(\frac{E_n^\pm(u, \phi)}{\Delta^*} \right) + \frac{L_g}{\xi_N(u)} \left(\frac{E_n^\pm(u, \phi)}{\Delta^*} \right) \pm \alpha(u, \phi) = 2\pi n, \quad (\text{B2})$$

where $\hbar$ is the reduced Planck's constant and ξ_N is the BCS coherence length, which can be calculated according to

$$\xi_N(u) = \frac{\hbar v_F(u)}{2\Delta^*}, \quad (\text{B3})$$

where $v_F(u) = \sqrt{2(E_F - E(u))/m^*}$ is the Fermi velocity, where, $m^* = 0.041m_e$ is the effective mass, where m_e is the electron mass, $E_F = (\pi \hbar^2/m^*)n_s$ is the Fermi energy, and $E(u) = u^2 \hbar^2/8m^*W^2$ is the subband energy obtained with the approximation for an infinitely deep quantum well. α in Eq. (B2) is the effective phase determined from [37]

$$\alpha(u, \phi) = \tau(u) \cos \phi + (1 - \tau(u)) \cos \left[\left(\frac{L_g - 2a}{\xi_N(u)} \right) \left(\frac{E_n^\pm(u, \phi)}{\Delta^*} \right) \right], \quad (\text{B4})$$

where $a = L_g/2$ is the position of a point impurity potential in the channel. Note that τ is also dependent on the

subband index by combining Eq. (A2) with the relation connecting the mean free path with the Fermi velocity

$$\lambda(u) = \frac{m^*}{q} \mu v_F(u), \quad (\text{B5})$$

assuming μ is independent of u . For all $E_n^\pm(u, \phi)$ satisfying Eq. (B2), the contribution to the supercurrent is [37]

$$I_n^\pm(u, \phi) = \frac{2q}{\hbar} \frac{\partial E_n(u, \phi)}{\partial \phi}, \quad (\text{B6})$$

and the total discrete current contribution I_{discrete} becomes

$$I_{\text{discrete}}(u, \phi) = \sum_{\pm n} I_n^\pm(u, \phi). \quad (\text{B7})$$

The contribution from the continuous spectrum starts to become important for long superconductor/normal conductor/superconductor junctions, i.e., when the gate length becomes comparable to the coherence length. If the electrons or holes forming the Andreev levels have energy $|E| > \Delta^*$, it is possible for them to escape from the junction into the continuous spectrum. These leaky states lead to a finite lifetime of the Andreev levels and reduce the total supercurrent [37]. The continuous contribution to the total supercurrent in the limit where $T \rightarrow 0 \text{ K}$ can be written as [37]

$$I_{\text{continuum}}(u, \phi) = \frac{2q}{h} \tau(u) \int_{-\infty}^{-\Delta^*} |u_0^2 - v_0^2| \left(\frac{1}{D(E, -\alpha)} - \frac{1}{D(E, \alpha)} \right) \frac{\sin \phi}{\sin \alpha} dE, \quad (\text{B8})$$

where u_0 and v_0 are the standard BCS coherence factors

$$\begin{cases} 2u_0^2 = 1 + \frac{\sqrt{E^2 - \Delta^*}}{E}, \\ 2v_0^2 = 1 - \frac{\sqrt{E^2 - \Delta^*}}{E}, \end{cases} \quad (\text{B9})$$

and the function $D(E, \phi)$ is expressed as [37]

$$D(E, \phi) = u_0^4 + v_0^4 - 2u_0^2 v_0^2 \cos \left[\left(\frac{E}{\Delta^*} \right) \left(\frac{L_g}{\xi_N} \right) + \phi \right]. \quad (\text{B10})$$

APPENDIX C: TEMPERATURE DEPENDENCE OF MAR PEAKS

Figure 6 shows the MAR-peak positions of order $n = \pm 3, \pm 4$ as a function of temperature T for $V_{\text{gs}} = -0.12 \text{ V}$. The measured data agree well with the model of an induced superconducting gap with $\Delta^* = 110 \mu\text{V}$ and $\gamma_B = 1.8$.

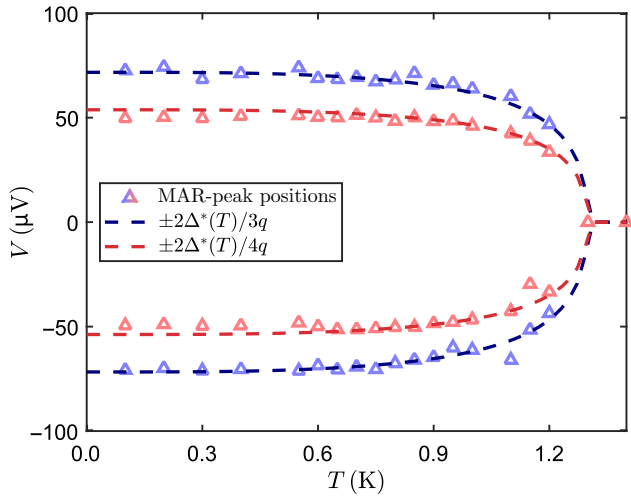


FIG. 6. MAR-peak positions for $n = \pm 3, \pm 4$ as a function of temperature T for $V_{gs} = -0.12$ V. Predictions of the temperature behavior of the induced gap are also shown.

- [1] K. Delfanazari, R. K. Puddy, P. Ma, T. Yi, M. Cao, Y. Gul, I. Farrer, D. A. Ritchie, H. J. Joyce, M. J. Kelly, and C. G. Smith, On-chip Andreev devices: hard superconducting gap and quantum transport in ballistic Nb-In_{0.75}Ga_{0.25}As-quantum-well-Nb Josephson junctions, *Adv. Mater.* **29**, 1701836 (2017).
- [2] L. Casparis, M. R. Connolly, M. Kjaergaard, N. J. Pearson, A. Kringhøj, T. W. Larsen, F. Kuemmeth, T. Wang, C. Thomas, S. Gronin, G. C. Gardner, M. J. Manfra, C. M. Marcus, and K. D. Petersson, Superconducting gatemon qubit based on a proximitized two-dimensional electron gas, *Nat. Nanotechnol.* **13**, 915 (2018).
- [3] A. Hertel, M. Eichinger, L. O. Andersen, D. M. T. Van Zanten, S. Kallatt, P. Scarlino, A. Kringhøj, J. M. Chavez-Garcia, G. C. Gardner, S. Gronin, M. J. Manfra, A. Gyenis, M. Kjaergaard, C. M. Marcus, and K. D. Petersson, Gate-tunable transmon using selective-area-grown superconductor-semiconductor hybrid structures on silicon, *Phys. Rev. Appl.* **18**, 034042 (2022).
- [4] J. I.-J. Wang, D. Rodan-Legrain, L. Bretheau, D. L. Campbell, B. Kannan, D. Kim, M. Kjaergaard, P. Krantz, G. O. Samach, F. Yan, J. L. Yoder, K. Watanabe, T. Taniguchi, T. P. Orlando, S. Gustavsson, P. Jarillo-Herrero, and W. D. Oliver, Coherent control of a hybrid superconducting circuit made with graphene-based van der Waals heterostructures, *Nat. Nanotechnol.* **14**, 120 (2019).
- [5] M. Hays, V. Fatemi, D. Bouman, J. Cerrillo, S. Diamond, K. Serniak, T. Connolly, P. Krogstrup, J. Nygård, A. Levy Yeyati, A. Geresdi, and M. H. Devoret, Coherent manipulation of an Andreev spin qubit, *Science* **373**, 430 (2021).
- [6] H. Riechert, S. Annabi, A. Peugeot, H. Duprez, M. Hantute, K. Watanabe, T. Taniguchi, E. Arrighi, J. Griesmar, J. D. Pillet, and L. Bretheau, The carbon nanotube gatemon qubit, *Nat. Commun.* **16**, 7197 (2025).
- [7] G. Butseraen, A. Ranadive, N. Aparicio, K. R. Amin, A. Juyal, M. Esposito, K. Watanabe, T. Taniguchi, N. Roch, F. Lefloch, and J. Renard, A gate-tunable graphene Josephson parametric amplifier, *Nat. Nanotechnol.* **17**, 1153 (2022).
- [8] D. Phan, P. Falthansl-Scheinecker, U. Mishra, W. M. Strickland, D. Langone, J. Shabani, and A. P. Higginbotham, Gate-tunable superconductor-semiconductor parametric amplifier, *Phys. Rev. Appl.* **19**, 064032 (2023).
- [9] J. Sarkar, K. V. Salunkhe, S. Mandal, S. Ghatak, A. H. Marchawala, I. Das, K. Watanabe, T. Taniguchi, R. Vijay, and M. M. Deshmukh, Quantum-noise-limited microwave amplification using a graphene Josephson junction, *Nat. Nanotechnol.* **17**, 1147 (2022).
- [10] E. D. Walsh, W. Jung, G.-H. Lee, D. K. Efetov, B.-I. Wu, K.-F. Huang, T. A. Ohki, T. Taniguchi, K. Watanabe, P. Kim, D. Englund, and K. Chung Fong, Josephson junction infrared single-photon detector, *Science* **372**, 409 (2021).
- [11] B. Huang, E. G. Arnault, W. Jung, C. Fried, B. J. Russell, K. Watanabe, T. Taniguchi, E. A. Henriksen, D. Englund, G.-H. Lee, and K. Chung Fong, Graphene calorimetric single-photon detector, arXiv:2410.22433.
- [12] G.-H. Lee, D. K. Efetov, W. Jung, L. Ranzani, E. D. Walsh, T. A. Ohki, T. Taniguchi, K. Watanabe, P. Kim, D. Englund, and K. Chung Fong, Graphene-based Josephson junction microwave bolometer, *Nature* **586**, 42 (2020).
- [13] R. Kokkonen, J.-P. Girard, D. Hazra, A. Laitinen, J. Govenius, R. E. Lake, I. Sallinen, V. Vesterinen, M. Partanen, J. Y. Tan, K. W. Chan, K. Y. Tan, P. Hakonen, and M. Möttönen, Bolometer operating at the threshold for circuit quantum electrodynamics, *Nature* **586**, 47 (2020).
- [14] J. Sarkar, K. Maji, A. Sunamudi, H. Agarwal, P. Samanta, A. Bhattacharjee, R. Rajkhowa, M. P. Patankar, K. Watanabe, T. Taniguchi, and M. M. Deshmukh, Kerr non-linearity enhances the response of a graphene Josephson bolometer, *Nat. Commun.* **16**, 7043 (2025).
- [15] T. Akazaki, H. Takayanagi, J. Nitta, and T. Enoki, A Josephson field effect transistor using an InAs-inserted-channel In_{0.52}Al_{0.48}As/In_{0.53}Ga_{0.47}As inverted modulation doped structure, *Appl. Phys. Lett.* **68**, 418 (1996).
- [16] C. T. Ke, C. M. Moehle, F. K. de Vries, C. Thomas, S. Metti, C. R. Guinn, R. Kallagher, M. Lodari, G. Scappucci, T. Wang, R. E. Diaz, G. C. Gardner, M. J. Manfra, and S. Goswami, Ballistic superconductivity and tunable π -junctions in InSb quantum wells, *Nat Commun.* **10**, 3764 (2019).
- [17] L. Olausson, P. Olausson, and E. Lind, Gate-controlled near-surface Josephson junctions, *Appl. Phys. Lett.* **124**, 042601 (2024).
- [18] C. Kittel, *Introduction to Solid State Physics* (John Wiley & Sons, Hoboken, NJ, 2005), Chap. 10, pp. 275.
- [19] P. Krogstrup, N. L. B. Ziino, W. Chang, S. M. Albrecht, M. H. Madsen, E. Johnson, J. Nygård, C. M. Marcus, and T. S. Jespersen, Epitaxy of semiconductor-superconductor nanowires, *Nat. Mater.* **14**, 400 (2015).
- [20] M. Kjaergaard, F. Nichele, J. Suominen, M. P. Nowak, M. Wimmer, A. R. Akhmerov, J. A. Folk, K. Flensberg, J. Shabani, C. J. Palmström, and C. M. Marcus, Quantized conductance doubling and hard gap in a two-dimensional semiconductor-superconductor heterostructure, *Nat. Commun.* **7**, 12841 (2016).
- [21] W. Mayer, J. Yuan, K. S. Wickramasinghe, T. Nguyen, M. C. Dartiaill, and J. Shabani, Superconducting proximity

- effect in epitaxial Al-InAs heterostructures, *Appl. Phys. Lett.* **114**, 103104 (2019).
- [22] M. Kjaergaard, H. J. Suominen, M. P. Nowak, A. R. Akhmerov, J. Shabani, C. J. Palmström, F. Nichele, and C. M. Marcus, Transparent semiconductor-superconductor interface and induced gap in an epitaxial heterostructure Josephson junction, *Phys. Rev. Appl.* **7**, 034029 (2017).
- [23] L. A. Glasser, Circuit models and applications of the superconducting field-effect transistor, *IEEE J. Solid-State Circuits* **24**, 1441 (1989).
- [24] F. Wen, J. Shabani, and E. Tutuc, Josephson junction field-effect transistors for Boolean logic cryogenic applications, *IEEE Trans. Electron Devices* **66**, 5367 (2019).
- [25] R. Rousset-Zenou, N. Aparicio, S. Messelot, R. D. Schlosser, M. Bjergfelt, J. Renard, M. Hocevar, and J. Nygård, Gate-tunable Josephson parametric amplifiers based on semiconductor nanowires, arXiv:2510.00305v1
- [26] W. Lu, A. Guo, A. Vardi, and J. A. del Alamo, A test structure to characterize nano-scale Ohmic contacts in III–V MOSFETs, *IEEE Electron Device Lett.* **35**, 178 (2014).
- [27] C. L. Hinkle, A. M. Sonnet, E. M. Vogel, S. McDonnell, G. J. Hughes, M. Milojevic, B. Lee, F. S. Aguirre-Tostado, K. J. Choi, H. C. Kim, J. Kim, and R. M. Wallace, GaAs interfacial self-cleaning by atomic layer deposition, *Appl. Phys. Lett.* **92**, 071901 (2008).
- [28] J. Ahn, T. Kent, E. Chagarov, K. Tang, A. C. Kummel, and P. C. McIntyre, Arsenic decapping and pre-atomic layer deposition trimethylaluminum passivation of Al₂O₃/InGaAs(100) interfaces, *Appl. Phys. Lett.* **103**, 071602 (2013).
- [29] F. Vigneau, R. Mizokuchi, D. C. Zanuz, X. Huang, S. Tan, R. Maurand, S. Frolov, A. Sammak, G. Scappucci, F. Lefloch, and S. De Franceschi, Germanium quantum-well Josephson FIELD-EFFECT TRANSISTORS AND INTERFEROMETERS, *Nano Lett.* **19**, 1023 (2019).
- [30] Y. Taur, D. A. Buchanan, W. Chen, D. J. Frank, K. E. Ismail, L. O. Shih-Hsien, G. A. Sai-Halas, R. G. Viswanathan, H.-J. C. Wann, S. J. Wind, and H. S. Wong, CMOS scaling into the nanometer regime, *Proc. IEEE* **85**, 486 (1997).
- [31] B. A. Aminov, A. A. Golubov, and M. Y. Kupriyanov, Quasiparticle current in ballistic constrictions with finite transparencies of interfaces, *Phys. Rev. B* **53**, 365 (1996).
- [32] G. E. Blonder, M. Tinkham, and T. M. Klapwijk, Transition from metallic to tunneling regimes in superconducting microconstrictions: Excess current, charge imbalance, and supercurrent conversion, *Phys. Rev. B* **25**, 4515 (1982).
- [33] M. Lundstrom and G. Jeong, *Near-Equilibrium Transport: Fundamentals and Applications* (World Scientific, Singapore, 2013), Chap. 3, pp. 47.
- [34] P. Olausson and E. Lind, Geometrical magnetoresistance as a tool for carrier mobility extraction in InGaAs MOSFETs, *IEEE Trans. Electron Devices* **70**, 5614 (2023).
- [35] L. Södergren, P. Olausson, and E. Lind, Cryogenic characteristics of InGaAs MOSFET, *IEEE Trans. Electron Devices* **70**, 1226 (2023).
- [36] K. K. Likharev, Superconducting weak links, *Rev. Mod. Phys.* **51**, 101 (1979).
- [37] P. F. Bagwell, Suppression of the Josephson current through a narrow, mesoscopic, semiconductor channel by a single impurity, *Phys. Rev. B* **46**, 12573 (1992).
- [38] A. Levchenko, A. Kamenev, and L. Glazman, Singular length dependence of critical current in superconductor/normal-metal/superconductor bridges, *Phys. Rev. B* **74**, 212509 (2006).
- [39] Y. Takane and K.-I. Imura, Quasiclassical theory of the Josephson effect in ballistic graphene junctions, *J. Phys. Soc. Jpn.* **81**, 094707 (2012).
- [40] M. Octavio, M. Tinkham, G. E. Blonder, and T. M. Klapwijk, Subharmonic energy-gap structure in superconducting constrictions, *Phys. Rev. B* **27**, 6739 (1983).
- [41] K. Flensberg, J. Bindslev Hansen, and M. Octavio, Subharmonic energy-gap structure in superconducting weak links, *Phys. Rev. B* **38**, 8707 (1988).
- [42] J. C. Cuevas, J. Hammer, J. Kopu, J. K. Viljas, and M. Eschrig, Proximity effect and multiple Andreev reflections in diffusive superconductor–normal-metal–superconductor junctions, *Phys. Rev. B* **73**, 184505 (2006).
- [43] X. Du, I. Skachko, and E. Y. Andrei, Josephson current and multiple Andreev reflections in graphene SNS junctions, *Phys. Rev. B* **77**, 184507 (2008).
- [44] Z. Wan, A. Kazakov, M. J. Manfra, L. N. Pfeiffer, K. W. West, and L. P. Rokhinson, Induced superconductivity in high-mobility two-dimensional electron gas in gallium arsenide heterostructures, *Nat. Commun.* **6**, 7426 (2015).
- [45] G. Niebler, G. Cuniberti, and T. Novotný, Analytical calculation of the excess current in the Octavio-Tinkham-Blonder-Klapwijk theory, *Supercond. Sci. Technol.* **22**, 085016 (2009).
- [46] S. N. Artemenko, A. F. Volkov, and A. V. Zaitsev, On the excess current in microbridges S-c-S and S-c-N, *Solid State Commun.* **30**, 771 (1979).
- [47] P. Samuelsson, A. A. Ingeman, G. Johansson, E. V. Bezuglyi, V. S. Shumeiko, G. Wendin, R. Kürsten, A. Richter, T. Matsuyama, and U. Merkt, Coherent current transport in wide ballistic Josephson junctions, *Phys. Rev. B* **70**, 212505 (2004).
- [48] R. Taboryski, T. Clausen, J. Bindslev Hansen, J. L. Skov, J. Kutchinsky, C. B. So/rensen, and P. E. Lindelof, Andreev reflections at interfaces between δ -doped GaAs and superconducting Al films, *Appl. Phys. Lett.* **69**, 656 (1996).
- [49] L. Olausson and E. Lind, Quasi-Ballistic InGaAs Josephson Field-Effect Transistor with Molybdenum Electrodes [Dataset]. Zenodo (2025), <https://doi.org/10.5281/zenodo.17515869>.
- [50] A. Khakifirooz, O. M. Nayfeh, and D. Antoniadis, A simple semiempirical short-channel MOSFET current–voltage model continuous across all regions of operation and employing only physical parameters, *IEEE Trans. Electron Devices* **56**, 1674 (2009).