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Anomalous metallic phase in tunable destructive superconductors

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Multiply connected superconductors smaller than the coherence length show destructive superconductivity, characterized by reentrant quantum phase transitions driven by magnetic flux. We investigate the dependence of destructive superconductivity on flux, transverse magnetic field, temperature, and current in InAs nanowires with a surrounding epitaxial Al shell, finding excellent agreement with mean-field theory across multiple reentrant transitions. Near the crossover between destructive and nondestructive regimes, an anomalous metal phase is observed with temperature-independent resistance, controlled over two orders of magnitude by a millitesla-scale transverse magnetic field.

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Quantum phase transitions [1,2] in conventional superconductors serve as prototypes for related effects in more complex, strongly correlated systems [3], including heavy-fermion materials [4] and high-temperature superconductors [5]. While low-temperature superconductors are well understood in bulk, new phenomena can arise in mesoscopic samples and reduced dimensionality [6,7]. For instance, in two-dimensional films, electrons theoretically condense into either a superconductor or insulator in the low-temperature limit [8]. Yet, in many instances, an anomalous metallic state with temperature-independent resistance is found at low temperatures [9]. In one-dimensional wires, incoherent phase slips can destroy superconductivity [10] or give rise to an anomalous metallic state [11], while coherent quantum phase slips can lead to superposition of quantum states enclosing different numbers of flux quanta [12], potentially useful as a qubit [13].

Multiply connected superconductors provide an even richer platform for investigating phase transitions. Fluxoid quantization in units of $\Phi_0 = \hbar/2e$ [14,15], reveals not only electron pairing but a complex macroscopic order parameter, $\Delta e^{ij}$ [6,16]. The same physical mechanism underlies the Little-Parks effect, a periodic modulation of the transition temperature, $T_C$, of a superconducting cylinder with magnetic flux period $\Phi_0$ [17]. For hollow superconducting cylinders with diameter, $d$, smaller than the coherence length, the modulation amplitude can exceed zero-field transition temperature, $T_{C0}$, leading to a reentrant destruction of superconductivity near odd half-integer multiples of $\Phi_0$ [18–20].

Early experimental investigation of the destructive Little-Parks effect reported reentrant superconductivity interrupted by an anomalous-resistance phase around applied flux $\Phi_0/2$ [21]. Subsequent experiments showed a low-temperature phase with normal-state resistance, $R_N$, around $\Phi_0/2$, but did not display fully recovered superconductivity at higher flux [22]. Several theoretical models were proposed to interpret these different scenarios [23–25], but no consensus emerged.

Here, we report a study of the Little-Parks effect across the transition from destructive to nondestructive regimes, in InAs nanowires with a thin epitaxial cylindrical Al shell. Remarkable agreement with Ginzburg-Landau mean-field theory is observed across multiple reentrant lobes as a function of flux, temperature, and current bias, using independently measured material and device parameters. We then investigate a field-tunable crossover from nondestructive to destructive regime. At the boundary, an anomalous metal phase is identified, characterized by a temperature-independent resistance that can be tuned over two orders of magnitude using small changes in perpendicular magnetic field, $B_L$. We interpret these results in terms of tunneling between adjacent fluxoid states with different phase winding numbers giving rise to an anomalous metallic phase. As noted previously [24], the appearance of a field-tunable temperature-independent resistance does not emerge naturally from simple models. The basic mechanism leading to a field-tunable saturating resistance remains mysterious.

The devices we investigated were made using InAs nanowire grown by the vapor-liquid-solid method using molecular beam epitaxy (MBE). Following wire growth, an

![FIG. 1. (a) Colorized material-sensitive scanning electron micrograph of InAs-Al hybrid nanowire cross section. The full wire diameter $d_B$, core diameter $d_C$, and shell thickness $t_S$ are indicated by dashed arrows. (b) Representative color-enhanced micrograph of a device (wire B) consisting of an InAs core (green) with Al shell (gray), contacted with Ti/Au leads (yellow). The device can be operated in voltage ($V$) and current ($I$) bias measurement setups.](image-url)
epitaxial Al layer was grown within the MBE chamber while rotating the sample stage, resulting in a full cylindrical Al shell coating the wire [26], as shown in Fig. 1(a). Subsequent fabrication used standard electron-beam lithography, deposition, etching, and liftoff, as described elsewhere [27]. Devices were operated in two configurations [Fig. 1(b)]: In the first configuration, four Au contacts were made to the Al shell allowing four-wire resistance measurements; in the second, an additional tunneling contact to the InAs core at the end of the wire was used as a tunnel probe, giving local density of states, as discussed in Ref. [27]. We investigated wires from three growth batches, denoted A, B, and C, with different core diameters, $d_C$, and shell thicknesses, $t_S$ (see Supplemental Material [28]). Transport measurements were carried out in a dilution refrigerator with a three-axis vector magnet and base temperature of 20 mK.

Carrier density in the InAs core is predominantly at the Al interface due to band bending [30,31]. Moreover, the density of carriers in Al is orders of magnitude higher than in InAs. We may therefore consider current to be carried by a density of carriers in Al is orders of magnitude higher than in InAs. We may therefore consider current to be carried by the parameter $\xi_A$, where $\xi = \frac{\alpha}{2\pi k_B T_C(\alpha)}$ is the applied flux quantum, as described by Abrikosov-Gorkov expression,

$$\ln \left( \frac{T_C(\alpha)}{T_{C0}} \right) = \Psi \left( \frac{1}{2} \right) - \Psi \left( \frac{1}{2} + \frac{\alpha}{2\pi k_B T_C(\alpha)} \right),$$

(1)

where $\Psi$ is the digamma function [32]. Following Refs. [20,22,23], the pair-breaking parameter for a hollow cylinder with wall thickness $t_S$ in a parallel magnetic field $B_\| \parallel$ is given by

$$\alpha = \frac{4 \xi_S^2 T_{C0}}{A_F} \left[ \left( \frac{n - \Phi}{\Phi_0} \right)^2 + \frac{t_S^2}{d_F^2} \left( \frac{\Phi^2}{\Phi_0^2} + \frac{n^2}{3} \right) \right],$$

(2)

where $\xi_S$ is the zero-field superconducting coherence length, $A_F$ is the area of the cylinder cross section, the integer $n$ is the fluxoid quantum number, $\Phi$ is the applied flux, and $d_F$ is the diameter of the cylinder [Fig. 1(a)]. Taking the dirty-limit expression for $\xi_S = \sqrt{\pi k_B T C_0 / 24 \Phi_0}$ with the Fermi velocity $v_F$ and mean free path $l_F$, we note that all parameters can either be measured directly from the micrograph of the device or from independent transport measurements (see Supplemental Material [28]).

Differential shell resistances, $R_S = dV_S/dI_S$, for wires A and B are shown in Fig. 2 as a function of $B_\|$, and temperature, $T$. Wires A and B have similar core diameters, $d_C \sim 135 \text{ nm}$, but different shell thicknesses. For wire A, with $t_S = 7 \text{ nm}$, $T_C$ is finite throughout the measured range, and varies periodically with applied axial flux with amplitude $\sim 0.4 \text{ K}$ with no clear envelope reduction up to $B_\| = 0.4 \text{ T}$. Normal-state resistance of the wire yields a coherence length $\xi_S = 70 \text{ nm}$, smaller than $d_C$ (see Supplemental Material [28]). In contrast, wire B, with $t_S = 25 \text{ nm}$, has $\xi_S = 180 \text{ nm} > d_C$, and shows destructive regimes around $\pm \Phi_0/2$ and $\pm 3\Phi_0/2$. Resistances in these destructive regimes remain equal to the normal-state resistance, $R_S = R_N$, to the lowest measured temperatures.

The absence (presence) of the destructive regime in wire A (B) is consistent with the criterion of the superconducting coherence length being smaller (larger) than the wire diameter [18]. To be more quantitative, we plot in Fig. 2 theoretical curves marking the superconductor-metal transition based on Eqs. (1) and (2) with independently measured wire parameters, using either the measured zero-field critical temperature, $T_{C0}$, or equivalently, the spectroscopically measured zero-field superconducting gap, $\Delta_0$ (Fig. S1 in the Supplemental Material [28]), which was consistent with the BCS relation $\Delta_0 = 1.76k_B T_{C0}$ [6]. Figure 2 demonstrates the remarkably good agreement found between experiment and theory. The observed increase of $T_C$ with decreasing $t_S$ is consistent with enhanced energy gaps for thin Al films [33].

Similar to the effects of flux-induced circumferential supercurrent, a dc current, $I_S$, applied along the wire can also drive the shell normal. The field-dependent critical current $I_{C}(\alpha)$ can be related to the corresponding critical temperature, $T_C(\alpha)$,

$$I_{C}(\alpha) = I_{C0} \left( \frac{T_C(\alpha)}{T_{C0}} \right)^{3/2},$$

(3)

where $I_{C0}$ is the zero-field critical current [34].

Base-temperature $I_S$-$B_\|$ phase diagrams for wires A and B are shown in Fig. 3. The data are taken sweeping from negative to positive bias, so show retrapping currents for $I_S < 0$ and switching current for $I_S > 0$, both of which are proportional to the critical current, $I_C$ [6]. Similar to the transition temperature, $I_C$ was observed to be $\Phi_0$ periodic in flux for both wires as expected from Eq. (3). For wire A, a
The transition of the wire C from being nondestructive at a long range. (b) Same as (a), but measured for wire B with shell thickness $\sim 13\Phi_0$, corresponding to a finite, but smaller than normal-state resistance. (c) Same as (a), but at $B_\perp = 12$ mT. Around half-flux quantum and zero-current bias, an anomalous phase develops with a finite, but smaller than normal-state resistance. (c) Same as (a), but measured at $B_\perp = 13$ mT. Around half-flux quantum $R_S$ remains at normal-state value even at $I_S = 0$. The theory curve in (d) is Eq. (3) computed with $\alpha_\perp$. Critical current evolution $I_C$ close to the crossover, around $B_\perp \sim 12$ mT, along with superimposed theory curves based on Eqs. (1)–(4). Note that unlike the situation far from the crossover [Fig. 5(a)], where theory and experiment agree well, in the vicinity of the crossover Figure 4 examines this resistive state close to the crossover, around $B_\perp \sim 12$ mT, along with superimposed theory curves based on Eqs. (1)–(4). Note that unlike the situation far from the crossover [Fig. 5(a)], where theory and experiment agree well, in the vicinity of the crossover [Figs. 5(b) and 5(c)] mean-field theory predicted $T_C$ deviates from the temperatures where the shell displays $R_N$. 

Theoretically, the effect of $B_\perp$ on the superconducting transition can be accounted for by introducing an additional pair-breaking term [3],

$$\alpha_\perp = \frac{4\xi_0^2 T_C \Phi_0^2}{\Phi_0^2},$$  

where $\Phi_0 = B_0 A_F$. Figure 4 shows the theoretical transitions based on Eqs. (1)–(4) using $\alpha = \alpha_\parallel + \alpha_\perp$ [35] superimposed on experimental data.
Temperature dependence of $R_S$ around $-\Phi_0/2$ for several values $B_\perp$ near the conventional-destructive crossover are shown in Fig. 5(d). Throughout this regime, $R_S$ was found to saturate to a temperature-independent value, which can be tuned over two orders of magnitude with small changes in $B_\perp$. In contrast, a $R_S$-$T$ trace taken close to the second destructive regime, not near a crossover ($B_\perp = 12$ mT and $B_\parallel = 52$ mT) remains temperature dependent down to the base temperature (Fig. S2 in the Supplemental Material [28]). Qualitatively similar anomalous $R_S$ saturation was also observed for different $B_\parallel$ values at a fixed $B_\perp$ (see Fig. S3 in the Supplemental Material [28]). At base temperature the evolution of $R_S$ as a function of $B_\perp$ shows a steplike increase that is mostly pronounced around $\pm \Phi_0/2$ [see Fig. 5(e)].

A possible explanation for the saturation of $R_S$ in terms of disorder-induced variations of $\Delta$, separating the shell into normal and superconducting segments [23], was tested by examining saturation effects in three segments of the same wire (Fig. S4 in the Supplemental Material [28]). It was found that all segments behaved the same, arguing against long-range variation in $\Delta$ on the scale of the separation of contacts. We also note that the anomalous resistance develops predominantly above the theoretical $T_c$, where the sample is expected to be in the normal state (Fig. S5 in the Supplemental Material [28]).

The steplike increase of $R_S$ with $B_\perp$ shown in Fig. 5(e) is reminiscent of phase slips, similar to the ones reported in Refs. [10,35], except here they are activated by perpendicular field rather than temperature. This suggests a picture in which anomalous saturating resistance results from quantum fluctuations not captured by mean-field theory. In general, the probability of a transverse phase slip across a weak link is proportional to $\exp(-R_2/R_N)$, with the resistance quantum $R_2$, and therefore is exponentially small for wire C [36]. However, near one-half-flux quantum, states with consecutive phase windings around the shell are degenerate, allowing quantum fluctuations to play a role. We note that both deep in the nondestructive regime [Fig. 2(a)] and deep in the destructive regime [Fig. 2(b)], no anomalous phase is observed.

Previous theoretical work [16,23] argued that the ratio of $d\Phi_0/2$ to $\lambda^2$ controls the order of the superconductor-metal transition. The present experiments span the range, with wires A and B having $d\Phi_0/2 < \lambda^2$, whereas wire C has $d\Phi_0/2 \gtrsim \lambda^2$. We have not observed systematic qualitative difference across this ratio. A detailed investigation of the order of the transition, and its affect on the anomalous metallic phase, would make an interesting future study.

In summary, we have investigated destructive and non-destructive Little-Parks effect in InAs nanowires fully
covered with epitaxial Al. Excellent agreement with Ginzburg-Landau mean-field theory was obtained across multiple reentrant quantum phase transitions using independently measured device and material parameters. Millitesla-scale perpendicular magnetic field was used to tune the crossover between destructive and nondestructive regimes, yielding an anomalous metallic phase around one-half-flux quantum with a temperature-independent resistance ranging over two orders of magnitude controlled by small changes in perpendicular field. This field-controllable anomalous phase is not explained by existing theory, but presumably involves quantum fluctuations between winding numbers of superconducting phase around the cylindrical superconducting shell.

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[1] S. L. Sondhi, S. M. Girvin, J. P. Carini, and D. Shahar, Rev. Mod. Phys. 69, 315 (1997).
[2] T. Vojta, Annu. Phys. (NY) 9, 403 (2000).
[3] N. Shah and A. Lopatin, Phys. Rev. B 76, 094511 (2007).
[4] Q. Si and F. Steglich, Science 329, 1161 (2010).
[5] M. R. Norman, Science 332, 196 (2011).
[6] M. Tinkham, Introduction to Superconductivity, 2nd ed. (McGraw-Hill, New York, 1996).
[7] A. Del Maestro, B. Rosenow, and S. Sachdev, Ann. Phys. (NY) 324, 523 (2009).
[8] A. M. Goldman, Int. J. Mod. Phys. B 24, 4081 (2010).
[9] A. Kapitulnik, S. A. Kivelson, and B. Spivak, Rev. Mod. Phys. 91, 011002 (2019).
[10] A. Bezryadin, C. N. Lau, and M. Tinkham, Nature (London) 404, 971 (2000).
[11] A. D. Zaikin, D. S. Golubev, A. van Otterlo, and G. T. Zimányi, Phys. Rev. Lett. 78, 1552 (1997).
[12] O. V. Astafiev, L. B. Ioffe, S. Kafanov, Yu. A. Pashkin, K. Yu. Arutyunov, D. Shahar, O. Cohen, and J. S. Tsai, Nature (London) 484, 355 (2012).
[13] J. E. Mooij and C. J. P. M. Harmans, New J. Phys. 7, 219 (2005).
[14] B. S. Deaver, Jr. and W. M. Fairbanks, Phys. Rev. Lett. 7, 43 (1961).
[15] H. Doll and M. Nåbauer, Phys. Rev. Lett. 7, 51 (1961).
[16] D. H. Douglass, Jr., Phys. Rev. 132, 513 (1963).
[17] W. A. Little and R. D. Parks, Phys. Rev. Lett. 9, 9 (1962).
[18] P.-G. de Gennes, C. R. Acad. Sci. Paris 292, 279 (1981).
[19] R. M. Arutyunyan and G. F. Zharkov, Zh. Eksp. Teor. Fiz. 79, 245 (1980) [Sov. Phys. JETP 52, 124 (1980)].
[20] G. Schwiete and Y. Oreg, Phys. Rev. Lett. 103, 037001 (2009).
[21] Y. Liu, Yu. Zadorozhny, M. M. Rosario, B. Y. Rock, P. T. Carrigan, and H. Wang, Science 294, 2332 (2001).
[22] I. Sternfeld, E. Levy, M. Eshkol, A. Tsuchikawa, M. Karpovski, H. Shtrikman, A. Kretinin, and A. Palevski, Phys. Rev. Lett. 107, 037001 (2011).
[23] V. H. Dao and L. F. Chibotaru, Phys. Rev. B 79, 134524 (2009).
[24] O. Vafek, M. R. Beasley, and S. A. Kivelson, arXiv:cond-mat/0505688.
[25] A. V. Lopatin, N. Shah, and V. M. Vinokur, Phys. Rev. Lett. 94, 037003 (2005).
[26] P. Krosgstrup, N. L. B. Ziino, W. Chang, S. M. Albrecht, M. H. Madsen, E. Johnson, J. Nygård, C. M. Marcus, and T. S. Jespersen, Nat. Mater. 14, 400 (2015).
[27] S. Vaitiekūnas, M.-T. Deng, P. Krosgstrup, and C. M. Marcus, arXiv:1809.05513.
[28] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevB.101.060507 for more detailed device description and additional measurements, see also Ref. [29].
[29] C. Kittel, Introduction to Solid State Physics, 8th ed. (Wiley, New York, 2005).
[30] A. E. Mikkelsen, P. Kotetes, P. Krosgstrup, and K. Flensberg, Phys. Rev. X 8, 031040 (2018).
[31] A. E. Antipov, A. Bargerbos, G. W. Winkler, B. Bauer, E. Rossi, and R. M. Lutchyn, Phys. Rev. X 8, 031041 (2018).
[32] A. A. Abrìkovsøv and L. P. Gor’kov, Zhur. Eksp. Teor. Fiz. 39, 1781 (1960) [Sov. Phys. JETP 12, 1243 (1961)].
[33] N. A. Court, A. J. Ferguson, and R. G. Clark, Supercord. Sci. Technol. 21, 015013 (2008).
[34] J. Bardeen, Rev. Mod. Phys. 34, 667 (1962).
[35] A. Rogachev, A. T. Bollinger, and A. Bezryadin, Phys. Rev. Lett. 94, 017004 (2005).
[36] M. Vanević and Y. V. Nazarov, Phys. Rev. Lett. 108, 187002 (2012).