in lower panels) lies in the magnetic equatorial plane. The reconnection potential is arbitrary up to a constant of integration, which was chosen to ensure that, for each run, the zero value of reconnection potential in the magnetotail is magnetically connected to the zero potential in the ionosphere (18).

The reconnection potentials versus MLT with offsets chosen in this way (Fig. 4) exhibit the same dusk-to-dawn ratios of extrema as the ionospheric potentials (to within 1% for the values in Fig. 3). The minimum (duskside) potential is larger in magnitude than maximum (dawnside) potential for the causal and auroral depletion runs and equal in magnitude (within numerical error) for the uniform run. The average reconnection rate is higher on the nightside than on the dayside and more spatially limited. As expected from the distribution of fast flows (Fig. 3), the reconnection rate, per unit length x-line, averaged over a 1-hour MLT segment to either side of or = 0 (Table 1), is larger in the premidnight sector, tailward of the fast flow, in the causal-empirical run and larger postmidnight in the Hall depletion run. The reconnection potential difference calculated as  =  -  exceeds the cross polar cap potential ( - ) in the ionosphere by 12, 6, and 14% for the uniform, causal, and auroral depletion runs, respectively (18).

Hall currents flow antiparallel to convection streamlines in the ionosphere (Fig. 3). If the Hall conductance develops a gradient parallel to a streamline—e.g., due to an enhancement in electron precipitation—the Hall current becomes discontinuous if other effects do not intervene. Its excess current could be diverted into a field-aligned current flowing through the magnetosphere to alleviate charge accumulation at the discontinuity, or a charge accumulation could polarize the ionospheric plasma and introduce a secondary electric field orthogonal to the primary convection electric field E (directed equatorward in a nightside auroral conductance band and sunward poleward of it). The resulting field, when added to the primary field, locally rotates the direction of convection and the Hall currents while driving secondary meridional Pedersen currents. This so-called Cowling effect produces the Harang reversal (24) in the nightside convection throat and increases the Joule dissipation in a zonally limited channel of enhanced conductance (25).

The Cowling effect dominates the electrodynamics of M-I coupling on the 100-km-scale ionospheric resolution of the global simulations. The sunward electric field in the polar cap produces a CW rotation of polar convection and a duskward drift in the magnetotail lobes. The secondary Pedersen current flowing equatorward in the nightside conductance band is supplied by a plasmasheet dynamo that generates field-aligned currents from a tailward current system (24). The resulting duskward, bulk Lorentz force presumably moves the otherwise symmetric plasmasheet flows toward the premidnight sector to produce the observed asymmetry in reconnection rates and plasmasheet fast flows.

These results demonstrate the intricate interplay between the SW-M-I interaction and ionospheric Hall conduction in regulating magnetotail reconnection.

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SUPPLEMENTARY MATERIALS

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Methods
Supplementary Text
Table S1
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SUPERCONDUCTIVITY

Fermi arcs in a doped pseudospin-1/2 Heisenberg antiferromagnet

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High-temperature superconductivity in cuprates arises from an electronic state that remains poorly understood. We report the observation of a related electronic state in a noncuprate material, strontium iodide (Sr3IrO4), in which the distinct cuprate fermiology is largely reproduced. Upon surface electron doping through in situ deposition of alkali-metal atoms, angle-resolved photoemission spectra of Sr3IrO4 display disconnected segments of zero-energy states, known as Fermi arcs, and a gap as large as 80 meV/100-electron volts. Its evolution toward a normal metal phase with a closed Fermi surface as a function of doping and temperature parallels that in the cuprates. Our result suggests that Sr3IrO4 is a useful model system for comparison to the cuprates.

Although the mechanism of high-temperature superconductivity (HTSC) remains an open question, it is commonly believed that certain distinct features of cuprates are essential to HTSC: spin-1/2 moment on a quasi-two-dimensional (2D) square lattice, Heisenberg antiferromagnetic coupling, and no orbital degeneracy. A minimal model based on this assumption can reproduce much of the phenomenology of the cuprates (1). Within this framework, it would be informative to realize the key features of cuprates in a different material (2). The 5d transition metal oxide Sr3IrO4, with a 5-valence shell, is a Mott insulator in which the orbital degeneracy is removed through strong spin-orbit coupling (3, 4). Despite strong entanglement of spin and orbital degrees of freedom, the resulting pseudospins (with |J| = 1/2 quantum number) exhibit the spin
that derive from carrier doping of Sr$_2$IrO$_4$ (7). As a result, despite the very different starting electronic structures, the effective low-energy physics of Sr$_2$IrO$_4$ can be described by the same minimal model as developed for the cuprates, with comparable magnitudes of coupling constants, suggesting possible unconventional HTSC upon carrier doping of this material (8). Indeed, a recent Monte Carlo study predicts a d-wave superconducting phase for the electron-doped case (9).

However, thus far, little is known experimentally about the nature of the electronic phases that derive from carrier doping of Sr$_2$IrO$_4$ (10, 11). One of the outstanding questions is whether the analogy to cuprates can be extended to the case of a half-filled band. $k_F$, Fermi momentum along the diagonal line.

Fig. 1. Fermiology evolution of Sr$_2$IrO$_4$ at $T = 70$ K. Constant-energy intensity maps taken at $E_F$ for surface coverage of (A) 0.5 ML, (B) 0.75 ML, (C) 0.85 ML, and (D) 1 ML, shown over a quadrant of the Brillouin zone of the undistorted square lattice. The actual Brillouin zone respecting the $\sqrt{2} \times \sqrt{2}$ superstructure of the lattice distortion coincides with the magnetic Brillouin zone whose boundary is indicated by the dashed lines. The intensity maps were normalized by the angular profile of the intensity along the FS measured at 1 ML. The map was symmetrized with respect to the diagonal line connecting (0, 0) and ($\pi$, $\pi$). $k_x$ and $k_y$ momenta along two Ir–O–Ir bonds in the IrO$_2$ plane. (E) Electron count based on the volume enclosed by the apparent large FS centered at the $\Gamma$ point and the distance from the node to ($\pi/2$, $\pi/2$) as a function of coverage. Zero of the left vertical axis corresponds to the case of a half-filled band.

dynamics of a Heisenberg antiferromagnet (5) with the nearest-neighbor magnetic exchange coupling having an energy scale on the order of 60 to 100 meV (6, 7). As a result, despite the very different starting electronic structures, the effective low-energy physics of Sr$_2$IrO$_4$ can be described by the same minimal model as developed for the cuprates, with comparable magnitudes of coupling constants, suggesting possible unconventional HTSC upon carrier doping of this material (8). Indeed, a recent Monte Carlo study predicts a d-wave superconducting phase for the electron-doped case (9).

However, thus far, little is known experimentally about the nature of the electronic phases that derive from carrier doping of Sr$_2$IrO$_4$ (10, 11). One of the outstanding questions is whether the analogy to cuprates can be extended to the case of a metallic phase, because Sr$_2$IrO$_4$ may evolve differently with doping. In particular, weakening of spin-orbit coupling may cause the system to revert to a multiband metal. On the other hand, if the single-band picture remains intact, one may expect the generic cuprate phenomenology to be reproduced in this material. To address this issue, we used an in situ surface doping technique, an approach that has been proven effective for surface doping control of complex oxides (12), to reveal the doping evolution of the electronic structures via the use of angle-resolved photoemission. By varying the surface coverage of potassium atoms deposited on the surface of the parent insulator Sr$_2$IrO$_4$, we follow the electronic evolution across the entire phase diagram from a Mott insulating state to a normal metallic state.
via a “strange" metallic phase that bears marked spectral similarities to those observed for the cuprates. We present our data in units of surface coverage rather than in electron doping concentration because of a highly nonlinear relation between the two.

Figure 1, A to D, shows the constant-energy intensity maps at the Fermi level ($E_F$) at $T = 70$ K for surface coverage ranging from 0.5 to 1 monolayer (ML). At 1 ML, we observe a closed, large Fermi surface (FS) indicative of a normal metal phase with the enclosed area of 54.3% of the 2D Brillouin zone (BZ) corresponding to 8.6% electron doping relative to a half-filled band. With decreasing coverage, however, we observe two gradual but conspicuous spectral changes that indicate formation of a distinct electronic phase. First, the intensity is much suppressed in an extended region near $(\pi, 0)$, breaking up the large FS into disconnected segments of Fermi arcs (J3, J4). Hereafter, we refer to the maximum (or minimum) intensity point as a “node” (or “antinode”) in analogy to cuprates. Second, the locus of the node shows a large shift toward $(\pi/2, \pi/2)$ (Fig. 1E). As a result, an attempt to interpret the arc as a part of the large FS encounters a qualitative inconsistency in the electron count based on Luttinger’s sum rule: The volume enclosed by the apparent large FS becomes less than half of the BZ for coverage below $\approx 0.6$ ML (Fig. 1E), in clear contradiction with electron doping. These results imply either a global change in the FS topology or an inconsistency with the notion of FS as the system departs from the normal metal phase observed at 1 ML. We note that an analogous qualitative inconsistency in the hole count occurs (I5) when interpreting the arcs in cuprates (I6) as part of an underlying large FS—an initial hint that the analogy to cuprates extends to the case of a metallic phase.

The Fermi arc metal phase can be characterized by an order parameter that depends on both the coverage and the temperature. Figure 2, A to D, shows the energy distribution curves (EDCs) for $T = 70$ K, symmetrized to remove the effect of the Fermi function. In the antinodal region, the gap starts to emerge below $T = 70$ K and 0.5 ML. A gap as large as $\approx 80$ meV, which rapidly decreases away from the antinode and becomes zero over a finite length of the contour that defines an arc (Fig. 2E). The gap at the antinode continuously diminishes with increasing surface coverage and becomes undetectable within the energy resolution ($\approx 18$ meV) beyond 0.85 ML (Fig. 2F).

Figure 3, A to C, shows the $E_F$ intensity map taken at $T = 30, 70$, and 110 K for surface coverage of 0.7 ML. Similar to the surface coverage dependence, with increasing temperature the arc expands its enclosed volume (when extrapolated to a large FS), crossing from less than half-filling at $T = 30$ K (48.9%) to more than half-filling (52.1% of BZ) at $T = 110$ K. Concomitantly, the arc elongates (Fig. 3F) with increasing temperature and evolves to a closed FS at $T = 110$ K. On cooling down to $T = 30$ K, the length of the arc shortens considerably, but a finite length of arc still remains (Fig. 3F), which is incompatible with a d-wave gap indicative of superconductivity. Unfortunately, measurement at lower temperatures to check for possible emergence of a d-wave gap was precluded by the high resistance of the sample, the bulk of which remains an insulator.

Consistent with the formation of a closed FS at high temperature, the antinode gap magnitude decreases with increasing temperature, as indicated by the gradual shift of the leading edge in the EDC (Fig. 3E and 3F, inset), in clear contrast to the temperature evolution at the node that shows only a minor thermal broadening (Fig. 3D). The gap starts to emerge below $T = 110$ K but grows slowly with decreasing temperature, in contrast to the typical behavior of an order parameter that sets in at a well-defined transition temperature. However, the substantial temperature dependence indicates a nontrivial origin and a many-body nature of the gap. We remark that the overall momentum (I9) and temperature (20, 21) dependences of the gap closely follow those observed for the pseudogap in cuprates.

All of the preceding results—carrier counting inconsistency and gap behavior—indicate that the parallel between Sr$_2$IrO$_4$ and cuprates persists in the metallic phase induced by electron doping and that the breaking up of a FS into disconnected segments is a general phenomenon of a system containing certain generic characteristics of cuprates. The formation of an arc, as opposed to “hot spots" at intersections with a magnetic zone boundary, indicates that the correlation strength in Sr$_2$IrO$_4$, which has been intensely debated (22–26), is in the intermediate-to-strong coupling regime (27). Thus, Mott physics and local correlations are essential for understanding of the physics of Sr$_2$IrO$_4$. Taking all results together, we conclude that the phenomenology of the Fermi arc must be accountable within the minimal description of the physics of doping a single-band, spin-1/2 antiferromagnetic Mott insulator, as shared by Sr$_2$IrO$_4$ and cuprates. Their differences in microscopic electronic structures help separate essential features of HTSC from material-specific features. In particular, whether the charge gap is of Mott-insulating (Sr$_2$IrO$_4$) or charge-transfer (cuprates) nature seems to have little bearing on the properties discussed here.

The nature of the electronic phase manifesting Fermi arcs in electron-doped Sr$_2$IrO$_4$ remains unclear at this point. A growing body of evidence that Fermi arcs are associated with distinct phases in cuprates that compete with HTSC—such as density waves (28), stripes (29, 30), and checkerboard orders (31)—prompts investigation of competing phases in carrier-doped Sr$_2$IrO$_4$. This, in turn, will help clarify which of these phases, if any, occur generically in proximity to HTSC. If,
SUPERCONDUCTIVITY

Observation of broken time-reversal symmetry in the heavy-fermion superconductor UPt₃

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Models of superconductivity in unconventional materials can be experimentally differentiated by the predictions they make for the symmetries of the superconducting order parameter. In the case of the heavy-fermion superconductor UPt₃, a key question is whether its multiple superconducting phases preserve or break time-reversal symmetry (TRS). We tested for asymmetry in the phase shift between left and right circularly polarized light reflected from a single crystal of UPt₃ at normal incidence and found that this so-called polar Kerr effect appears only below the lower of the two zero-field superconducting transition temperatures. Our results provide evidence for broken TRS in the low-temperature superconducting phase of UPt₃, implying a complex two-component order parameter for superconductivity in this system.

he heavy-fermion metal UPt₃ (1) is one of only a handful of unconventional superconductors (2, 3) exhibiting multiple superconducting phases (4–6). In the normal state, strong hybridization between itinerant platinum 5d electrons and localized uranium 5f moments results in an effective mass that is ~50 times that of free electrons (2). Below the Néel temperature $T_N \approx 5$ K, the local U moments order antiferromagnetically in the a-b plane (7). In zero magnetic field, two peaks in the specific heat at $T_c \approx 550$ mK and $T_c \approx 480$ mK indicate the presence of two superconducting states of differing symmetry, called the A and B phases, respectively (4–6). Pressure studies suggest that these two phases couple to, and are stabilized by, the antiferromagnetic order parameter (8). In finite magnetic fields, three distinct vortex phases are also observed (9–11). The phase diagram of UPt₃ therefore presents a particular challenge for models of unconventional superconductivity.

In the absence of a detailed understanding of the microscopic origins of this system, theoretical and experimental efforts center on identifying the structure of the macroscopic superconducting order parameter—the pair wavefunction, or gap. In the case of UPt₃, acceptable candidate order parameters should respect the $D_{5h}$ point-group symmetry of the underlying crystal lattice and should therefore transform under one or more representations of this group.

In this framework, many—but not all—experimentally studied of the superconducting states (1) favor an $E_u$ odd-parity triplet representation in which the gap is given by

$$\Delta(k_F) = \hat{\epsilon} |\eta_1(T)(k_x^2 - k_y^2)\hat{b}_z \pm 2i \cdot \eta_2(T)\hat{k}_x\hat{k}_y\hat{k}_z|$$

in a coordinate system where $\hat{\epsilon} \parallel \hat{c}$. Here, $\eta_1(T) = \sqrt{1 - (T/T_c)}$ is the real component of the superconducting order parameter, marking the onset of the A phase at $T_c$, whereas $\eta_2(T) = \sqrt{1 - (T/T_c)}$ introduces an additional imaginary component in the B phase at $T_c$. An order parameter of this form first breaks gauge symmetry in the A phase, in which it also exhibits fourfold rotational symmetry in the a-b plane distinct from the hexagonal symmetry of the crystal lattice. In the B phase, the order parameter becomes isotropic in the a-b plane as $T \rightarrow 0$ K, and the phase difference between the real and imaginary components imprints an overall angular momentum to the pair wave function. Hence, time-reversal symmetry (TRS) is broken in this phase, with the sign of the imaginary component determining the orientation (chirality) of the internal angular momentum of the pair along $\hat{z}$.

The results of Josephson interferometry experiments (12, 13) are consistent with the spatial symmetries of the $E_u$ order parameter of Eq. 1. However, attempts to observe TRS-breaking (TRSB) in UPt₃ via muon spin relaxation measurements have yielded conflicting results (15, 16). Moreover, recent thermal conductivity data (17) have been interpreted to support a gap function belonging to an $E_u$ representation that precludes TRSB in the B phase. Thus, the unresolved question of whether TRS is indeed broken in the B phase has become critical to determining the symmetry and hence the proper classification of the superconducting order parameter of UPt₃.

A general consequence of a TRSB order parameter (with a net moment oriented along the c