Heat transport study of field-tuned quantum criticality in $CeIrIn_5$

H. Shakeripour,[1,2,*] M. A. Tanatar,[2,†] C. Petrovic,[3,4] and Louis Taillefer[2,4,‡]
[1]*Department of Physics, Isfahan University of Technology, Isfahan 84156-83111, Iran*
[2]*Département de Physique and RQMP, Université de Sherbrooke, Sherbrooke, Canada*
[3]*Department of Physics, Brookhaven National Laboratory, Upton, New York 11973, USA*
[4]*Canadian Institute for Advanced Research, Toronto, Ontario, Canada*
(Received 7 October 2015; published 8 February 2016)

The in-plane electrical resistivity, ρ, and thermal conductivity, κ, of the heavy-fermion superconductor $CeIrIn_5$ were measured down to 40 mK in magnetic fields up to 11 T applied along the c axis. For all fields above $H_{c2} = 4$ T of filamentary superconductivity, we find that the ratio of heat and charge conductivities in the $T \to 0$ limit obeys the Wiedemann-Franz law, $\kappa/T = L_0/\rho$, where $L_0 = 2.45 \times 10^{-8}$ W Ω K^{-2} is the Sommerfeld value of the Lorenz number. The temperature-dependent parts of both the electrical and thermal resistivity, $w \equiv T/L_0\kappa$, follow the functional dependence expected for the Fermi liquid theory of metals with $\rho - \rho_0 = AT^2$, $w - w_0 = BT^2$, with $\rho_0 = w_0$ and $B \approx 2A$. The coefficient B does not show a significant field dependence even upon approaching $H_{c2} = 0.4$ T of the bulk superconducting state. The weak response to the magnetic field is in stark contrast with the behavior found in the closely related $CeCoIn_5$, in which the field-tuned quantum critical point coincides with H_{c2}. The value of the electron-electron mass enhancement, as judged by the A and B coefficients, is about one order of magnitude reduced in $CeIrIn_5$ as compared to $CeCoIn_5$ (in spite of the fact that the zero field γ_0 in $CeIrIn_5$ is twice as large as γ_0 in $CeCoIn_5$), which suggests that the material is significantly farther away from the magnetic quantum critical point at bulk H_{c2} and at all of the studied fields. A suppressed Kadowaki-Woods ratio in $CeIrIn_5$ compared to $CeCoIn_5$ suggests a notably more localized nature of f electrons in the compound.

DOI: 10.1103/PhysRevB.93.075116

I. INTRODUCTION

Ce-based heavy-fermion compounds $CeMIn_5$ (M = Co,Rh,Ir), known as the 115 family, show a variety of ground states at ambient pressure. $CeIrIn_5$ and $CeCoIn_5$ are bulk superconductors with the transition temperatures $T_c = 0.4$ and 2.3 K, respectively [1,2]. Filamentary superconductivity is observed in $CeIrIn_5$ below $T_{cf} \sim 1.2$ K [1]. $CeRhIn_5$, an antiferromagnet with Neél temperature $T_N = 3.8$ K, exhibits ambient pressure superconductivity at very low temperatures, $T_c \approx 0.1$ K [3,4], and with the same value of $T_c \approx 2.5$ K as $CeCoIn_5$ under pressure of 2 GPa [5].

Magnetic order can be induced in the superconducting members by Cd and Hg substitutions of In [6], so that phase diagrams with a rich interplay of magnetism, quantum critical points (QCPs) at $T_N \to 0$ [7], and superconductivity are found as functions of pressure and composition (see Refs. [8,9] for reviews). Doping studies found that $CeIrIn_5$ [10–12], similar to $CeCoIn_5$ [6,13], also lies in the vicinity of the magnetic quantum critical point, and thus its superconductivity is most likely magnetically mediated. Our studies of thermal conductivity in both $CeCoIn_5$ [14] and $CeIrIn_5$ [15,16], though, found significant deviations from the simple d-wave scenario, supported by many experiments [17–20].

A magnetic field can also be used as the nonthermal tuning parameter for the quantum critical point, and in $CeCoIn_5$, the field-tuned QCP coincides mysteriously with the superconducting upper critical field for magnetic-field orientations along both the c axis [21,22] and the ab plane [23] of the tetragonal crystals. The existence of the QCP is reflected in anomalous temperature dependences of electrical resistivity [2,13,24] (closely following a T-linear dependence) and of the specific heat (showing logarithmic divergence), in band-dependent divergence of cyclotron effective masses [25,26] and directional violation of the Wiedemann-Franz law [24,27].

In $CeRhIn_5$, field-induced magnetic order is found in the pressure range below $p_{max} = 2.5$ GPa, corresponding to maximum superconducting T_c [28]. Interestingly, similar phenomenology of QCP coinciding with the first-order transition at H_{c2} is found in closely related heavy-fermion Ce_2PdIn_8 [29], as well as in some other unconventional superconductors, including cuprates [30] and iron-based material KFe_2As_2 [31].

Of great interest is the question of whether field-tuned quantum criticality can be found in $CeIrIn_5$. The two materials, $CeCoIn_5$ and $CeIrIn_5$, have identical Fermi surfaces [32,33] and show very similar normal state properties [33], so it is natural to expect similarity in response to magnetic field. Despite this apparent similarity, heat capacity measurements [34–36] do not indicate critical behavior near the bulk upper critical field, finding nearly temperature-independent C/T at the lowest temperatures. Instead, heat capacity, torque, and resistivity studies suggest that $CeIrIn_5$ shows a metamagnetic transition with a critical field of 25 T, far above superconducting H_{c2} [36–38]. These studies, though, reveal several unusual features. While the electronic term $\gamma_0 \equiv C/T$ in heat capacity saturates at a constant value below $T_{FL} \approx 0.8$ K in a magnetic field of 1 T and remains nearly constant with magnetic field $H \parallel c$ in the range 1–17 T, with $\gamma_0 \approx 0.7$–0.8 J/mol K^2 [35,36], deviations from this value at high fields start at progressively lower T_{FL} [36]. Another study suggested that the superconducting dome of $CeIrIn_5$ is located inside the dome of the precursor phase, similar to the pseudogap state in the high-T_c cuprates [39,40], suppressed at a field of 6 T.

*hshakeri@cc.iut.ac.ir
†Current address: Ames Laboratory, Ames, Iowa 50011, USA.
‡Louis.Taillefer@USherbrooke.ca

We should keep in mind, though, that resistivity studies in $CeIrIn_5$ cannot probe the region close to bulk H_{c2}, hidden inside the domain of filamentary superconductivity [1,41]. Subtraction of the Schottky anomaly in heat capacity measurements always leaves some uncertainty in the behavior at the lowest temperatures. To gain additional insight into the response of $CeIrIn_5$ to application of a magnetic field, here we report the field evolution of the electronic transport properties of $CeIrIn_5$ using simultaneous electrical resistivity and thermal conductivity measurements, with the latter enabling us to probe the range of filamentary superconductivity close to bulk H_{c2}. We find that the Wiedemann-Franz law is universally obeyed in the $T \to 0$ limit for all fields where resistivity measurements are not affected by filamentary superconductivity, proving the true bulk nature of both electrical and heat transport. For all fields above bulk H_{c2}, we observe a Fermi-liquid (FL) temperature dependence of thermal resistivity, $w - w_0 = BT^2$, slightly decreasing with the field B coefficient and nearly field-independent T_{FL}. Notably smaller values of B in $CeIrIn_5$ when compared to $CeCoIn_5$, as well as their very mild dependence on magnetic field, suggest that $CeIrIn_5$ is significantly farther away from magnetic instability at H_{c2}, and the field-tuned criticality is significantly different in the two materials. We link this difference to the anomalously low value of the Kadowaki-Woods ratio of the T^2 coefficient in electrical resistivity versus the γ coefficient of the specific heat. This suggests a more localized character of magnetism in $CeIrIn_5$ than in $CeCoIn_5$.

II. EXPERIMENT

Single crystals of $CeIrIn_5$ were grown using the self-flux method [1]. As-grown crystals had a thin platelike shape, with large surfaces corresponding to the (001) basal plane of the tetragonal lattice. The two samples used in this study were from two different batches, but they revealed nearly identical results. Samples were cut into a rectangular shape with dimensions of $\sim 4 \times 0.1 \times 0.045$ mm^3 and $\sim 2 \times 0.1 \times 0.01$ mm^3. Four contacts to the samples were made by soldering 50 μm silver wires with indium solder, which resulted in a contact resistance typically $\sim$1 mΩ at low temperature. The same contacts were used to measure electrical resistivity and thermal conductivity in a dilution refrigerator, thus essentially eliminating the uncertainty of the geometric factor determination. The thermal conductivity was measured using a standard four-probe steady-state method with two RuO_2 chip thermometers calibrated *in situ* against a reference Ge thermometer. The electrical resistivity was measured using an LR700 ac resistance bridge, operating at a frequency of 16 Hz, by applying 0.1 mA excitation currents. To enable low-noise measurements of low resistance samples in identical conditions with thermal conductivity studies, measurements were performed in a separate dilution refrigerator run, during which the high-resistance wire coils of the thermal conductivity apparatus were temporary shorted with silver wires; see Ref. [24] for details.

The phonon contribution to thermal conductivity, κ_g, was determined from electrical resistivity and thermal conductivity measurements in heavily disordered samples of $Ce_{1-x}La_xIrIn_5$ with $x = 0.2$ and residual resistivity $\sim$10 $\mu\Omega$ cm; see Ref. [42] for details. In these heavily disordered samples, the temperature-dependent inelastic scattering contribution to resistivity can be neglected, and the electronic contribution can be directly determined from the Wiedemann-Franz law as $\kappa_e/T = L_0/\rho$, which gives $\kappa_g = \kappa - \kappa_e$. This contribution represents a minor correction in the temperature range below 1 K, which is the focus of this study.

III. RESULTS

In Fig. 1, we plot the temperature-dependent resistivity of $CeIrIn_5$ measured in magnetic fields of 0, 3.15, 3.6, 4, 8, and 11 T applied along the tetragonal c axis (top-left panel). In zero field, the filamentary superconductivity is observed below $T_{cf} = 1.2$ K, and the resistive transition is completely suppressed in a field of 4 T. The filamentary superconductivity

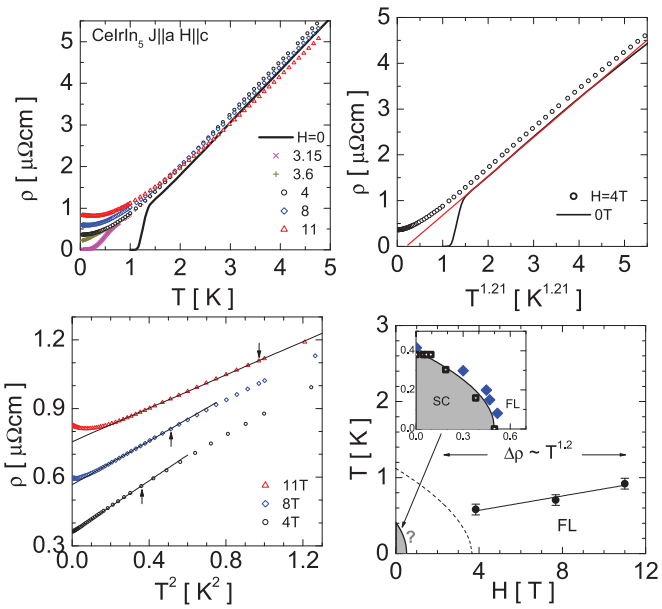

FIG. 1. Temperature-dependent in-plane electrical resistivity of $CeIrIn_5$ measured in magnetic fields of 0, 3.15, 3.6, 4, 8, and 11 T applied along the tetragonal c axis (top-left panel). The filamentary superconductivity, observed in zero field below $T_{cf} = 1.2$ K, is completely suppressed by $H_{c2f} \approx 4$ T, revealing a saturating $\rho(T)$ dependence on $T \to 0$. The top-right panel shows the data in zero field and in a field of 4 T plotted vs T^n with $n = 1.21$, best fit to the power-law function above the filamentary superconducting transition at $T_{cf} = 1.2$ K. The bottom-left panel shows $\rho(T)$ data plotted vs T^2 for magnetic field $H > H_{c2f}$. Linear dependence, found at low temperatures for all fields, suggests the validity of the Fermi-liquid picture predicting $\rho = \rho_0 + AT^2$. Arrows indicate the high-temperature end of the interval where this linear dependence is obeyed, marking a characteristic temperature T_{FL}. Note the low values of $T_{FL} \approx 0.7$ –1 K for all fields and the slowly decreasing slope of the $\rho(T^2)$ curves with magnetic field, suggesting a decreasing A coefficient (and thus effective mass) with increasing magnetic field. The bottom-right panel summarizes the phase diagram as found from resistivity, thermal conductivity (black squares) [15], and heat capacity (blue diamonds) [35] measurements. The dashed line in the main panel denotes the domain of filamentary superconductivity. We also show the field evolution of the Fermi-liquid properties as found from resistivity measurements (black dots) above H_{c2f}.

was shown to have the same angular dependence of H_{c2} as the bulk one [33,41], and it was suggested to originate from strain in the samples that was likely induced by the inclusion of a different phase [43]. A similar phenomenon is observed in Sr_2RuO_4 in which superconductivity with higher T_c is induced by the inclusion of Ru metal [44] and in some other materials including chain Nb_2Se_3 [45].

The $\rho(T)$ dependences at temperatures above 1 K are very similar for all magnetic fields. They clearly deviate from the functional form expected in Fermi-liquid theory, $\rho = \rho_0 + AT^2$. The residual electrical resistivity, ρ_0, of the samples was determined from measurements down to 40 mK in a magnetic field of 4 T, which was sufficient to completely suppress filamentary superconductivity. It was about $\sim$0.3 $\mu\Omega$ cm, showing the high purity of the samples. This value is notably lower than $\sim$0.5 $\mu\Omega$ cm in previous studies by Nair *et al.* [39] and $\sim$1.6 $\mu\Omega$ cm in a magnetic field of 12 T from Capan *et al.* [36]. The top-right panel shows the data in zero field and at 4 T plotted versus T^n with $n = 1.21$, which is the best fit to the power-law function for $T_{cf} < T < 3.5$ K. This value is very close to the previously reported $n = 1.27$ [10]. A similar power-law dependence with $n = 1.21$ is found for fields up to 11 T, with the range of power-law dependence continuing down to $T \approx 1$ K. At lower temperatures, all $\rho(T)$ curves show clear signs of saturation when T tends to zero, and at the lowest temperatures they follow the expectations of Fermi-liquid behavior. This is illustrated in the bottom-left panel of Fig. 1, in which resistivity data are plotted versus T^2, linearizing the dependence with the slope proportional to the A coefficient. Slight upward deviations at the lowest temperatures for 8 and 11 T curves are similar to those observed in $CeCoIn_5$ [22], and they are due to orbital magnetoresistance. For all magnetic fields $H > H_{c2f}$, the $\rho(T^2)$ plots reveal a range of linear dependence at the lowest temperatures. Deviations from the T^2 dependence at high temperatures mark a characteristic temperature T_{FL}, denoted by arrows. Note that T_{FL} has values in the range 0.7–1 K, with resistivity crossing over to T^n at slightly higher temperatures. T_{FL}, as determined from resistivity measurements, increases slightly with magnetic field, as shown in the phase diagram (bottom-right panel of Fig. 1). Due to a crossover character of deviations, definitions of T_{FL} have somewhat high errors, so it is difficult to make a reliable extrapolation of $T_{FL}(H)$ toward bulk $H_{c2} \approx 0.49$ T, determined from both thermal conductivity measurements [15,46] and heat capacity measurements [35]. The line in the bottom-right panel of Fig. 1 suggests extrapolation of T_{FL} to finite values in the $H = 0$ limit, suggesting the nonexistence of the field-tuned quantum critical point at H_{c2} or even in zero field. However, significant error bars of the T_{FL} determination make this extrapolation from field 4 T unreliable toward bulk H_{c2}.

To gain additional insight into the magnetic-field-induced evolution of the normal state in $CeIrIn_5$, we performed thermal conductivity measurements. These measurements are not affected by filamentary superconductivity, and they can be performed in the very proximity of bulk H_{c2}. In Fig. 2, we show the thermal conductivity of $CeIrIn_5$ measured in three magnetic fields: 0.52 T (slightly above the bulk $H_{c2} = 0.49$ T), 4 T (above H_{c2f}), and 10.7 T (far away from superconductivity). As can be seen from the 0.52 T curve, the κ/T versus T plot does not show any features upon entering the filamentary

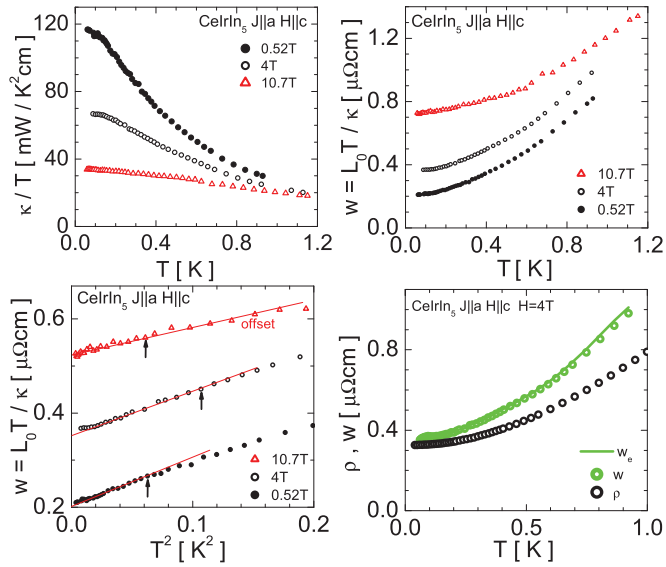

FIG. 2. Temperature-dependent in-plane thermal conductivity of $CeIrIn_5$, plotted as κ/T vs T, measured in magnetic fields of 0.52 T (slightly above bulk $H_{c2} = 0.49$ T), 4 T, and 10.7 T applied along the tetragonal c axis (top-left panel). The filamentary superconductivity does not affect the thermal conductivity of the samples, but it leads to an extrinsic violation of the Wiedemann-Franz law. Upon complete suppression of filamentary superconductivity in a magnetic field 4 T, the thermal resistivity $w \equiv L_0T/\kappa$ closely matches electrical resistivity ρ in the $T = 0$ limit, as shown in the bottom-right panel, clearly satisfying the Wiedemann-Franz law. Moreover, in the temperature range $T < 1$ K, in which the phonon contribution is negligible, the thermal resistivity w_e follows the same Fermi-liquid temperature dependence $w = w_0 + BT^2$. This dependence is closely obeyed for all magnetic fields studied (bottom panels), even in the vicinity of bulk H_{c2}. Due to larger noise in the data, the determination of T_{FL} in $w(T^2)$ plots is even less precise than that from the resistivity data (bottom-right panel).

zero-resistance state below approximately 1.0 K (not shown), leading to artificial violation of the Wiedemann-Franz law. As long as filamentary superconductivity is suppressed, thermal conductivity and electrical resistivity perfectly obey the Wiedemann-Franz law in the $T \to 0$ limit, with $\kappa/T\sigma = L_0$, where electrical conductivity $\sigma = 1/\rho$, and $L_0 = \frac{\pi^2}{3}(\frac{k_B}{e}^2) = 2.45 \times 10^{-8}$ W Ω K^{-2} is the Sommerfeld value of the Lorenz number, expected for fermionic systems. The validity of this ratio in the $T \to 0$ limit can be seen by direct comparison of the temperature-dependent thermal analog of electrical resistivity, $w \equiv L_0T/\kappa$, and resistivity (bottom-right panel in Fig. 2), which meet at $T \to 0$. $CeIrIn_5$ is a good metal, and therefore the phonon contribution to thermal conductivity can be completely neglected in the temperature interval of interest, $T < 1$ K. In Fig. 2 (bottom-right panel), we show this explicitly by subtracting the phonon contribution κ_g, as determined from measurements in disordered samples, from measured κ of pure $CeIrIn_5$. The resultant electronic part of thermal resistivity w_e is plotted in the bottom-right panel of Fig. 2 by a solid line. Close agreement is observed between w and w_e below 1 K, validating our conclusion about the negligible phonon correction.

IV. DISCUSSION

A. Temperature-dependent resistivity

As can be seen in the top panels of Fig. 1, the $\rho(T)$ curves in $CeIrIn_5$ strongly deviate from expectations of the Fermi-liquid functional dependence above $T \sim 1$ K, and they are well described by a $\Delta\rho \propto T^n$ dependence with $n = 1.21$. Observation of a very strong $\rho(T)$ dependence at such low temperatures [where phonon scattering becomes negligible but the strength of inelastic scattering, $\Delta\rho_{in} = \rho(1\ \text{K}) - \rho(0) = 0.6\ \mu\Omega\,\text{cm}$, is significantly larger than $\rho(0) = 0.3\ \mu\Omega\,\text{cm}$] suggests strong magnetic scattering in the compound, similar to $CeCoIn_5$ [2,33]. The power-law non-Fermi-liquid temperature-dependent resistivity with similar exponents n is not unusual in the heavy-fermion systems close to quantum critical points. For example, similar values of the power-law exponents are found in $CePd_2Si_2$ ($n \sim 1.2$) [7]. The low exponent $n = 1.21$ can be explained within spin fluctuation theory [47,48] for nearly antiferromagnetic metals in the framework of a spin-density-wave (SDW) scenario, though significantly lower $\Delta\rho_{in}/\rho_0$ are usually expected in this case [48]. In the three-dimensional SDW scenario, the exponents are usually expected to be close to $n = 1.5$ as found in three-dimensional (3D) $CeIn_3$ ($n \sim 1.5$) [7]. For the case of two-dimensional magnetic fluctuations, the exponents are usually expected to be close to $n = 1$ [49], as found in $YbRh_2Si_2$ [50,51] and $CeCu_{5.2}Ag_{0.8}$ ($n \sim 1$) [52]. In $CeCoIn_5$, the $\rho(T)$ in zero field is close to T-linear above T_c [2], but it becomes $n = 1.5$ at field-tuned QCP [27]. In this compound, the exponent n reveals significant anisotropy with $n = 1$ in $\rho_c(T)$ [24,53].

In $CeIrIn_5$, the same power-law behavior of resistivity up to 3.5 K is found for all magnetic fields. This observation suggests that despite the magnetic character of scattering, dominant scattering is not tuned by a magnetic field. This is in stark contrast with $CeCoIn_5$, which shows different power-law resistivity for low and high magnetic fields in the normal state [22,27].

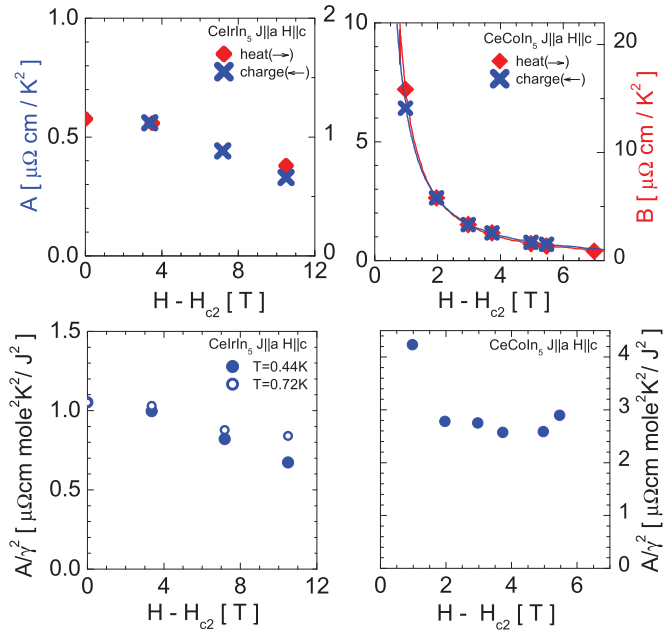

FIG. 3. Magnetic-field dependence of the T^2 coefficients of the electrical resistivity, A (blue crosses, left axes), and of the thermal resistivity, B (red diamonds, right axes), for $CeIrIn_5$ (top-left panel) and $CeCoIn_5$ (top-right panel). The data are plotted vs $H - H_{c2}$, which is the distance to a quantum critical point in $CeCoIn_5$. Note the tenfold difference of the A/B scales for the two cases. The field dependence $B(H)$ in $CeIrIn_5$ suggests no divergence of the effective mass toward bulk H_{c2}. Significantly smaller values of A and B coefficients compared to those in $CeCoIn_5$ suggest that $CeIrIn_5$ is significantly farther away from the quantum critical point in the phase diagram. A decrease of the A and B coefficients with magnetic field suggests moving away from the QCP. The bottom panels show Kadowaki-Woods ratios, A/γ^2, calculated based on the measured A values and the literature values of the specific heat [36]. Note that the heat capacity data show a notable temperature dependence at 11 T, as reflected in the higher ratio value at 0.72 K.

B. Field evolution of the effective mass

The temperature-dependent part of electronic thermal resistivity obeys the same functional dependence as electrical resistivity. However, because of the different Fermi surface integrals of heat and charge scatterings, the T^2 coefficients A for resistivity and B for thermal resistivity are not the same. Typically, for Fermi surfaces without sharp corners, the two coefficient are different by a factor of about 2 [42]. In the left panel of Fig. 3, we plot the magnetic-field dependence of A and B coefficients of the T^2 terms in the in-plane electrical (left axis) and thermal (right axis) resistivity of $CeIrIn_5$. The A and B coefficients in $CeIrIn_5$ are indeed close to obeying the $B = 2A$ relation, similar to $CeCoIn_5$ and most other cases [54–56]. The A and B coefficients in $CeIrIn_5$ closely follow each other as functions of magnetic fields, with B decreasing from approximately 1.2 to 0.8 upon increasing the field from 0.52 to 11 T. For reference, in the right panel we show similar data for $CeCoIn_5$ [27]. Note the tenfold difference in the scales for the two compounds. In sharp contrast to $CeCoIn_5$, the A and B coefficients in $CeIrIn_5$ do not reveal any tendency for divergence as functions of magnetic fields. The nearly tenfold difference in the A and B coefficients between $CeCoIn_5$ and $CeIrIn_5$ suggest a notable difference in the effective masses, reflecting close proximity to the quantum critical point in the case of $CeCoIn_5$ and a relatively large distance from it for $CeIrIn_5$.

This very weak field dependence of the A and B coefficients in $CeIrIn_5$ is consistent with specific heat [34,36], finding a weak field dependence of the electronic specific-heat coefficient γ_0, and the dHvA experiments [57] and cyclotron effective mass, m^* (for a field between 6 and 17 T) at low temperature [36,57]. In $CeCoIn_5$, the A coefficient changes from 7.5 to 0.5 $\mu\Omega\,\text{cmK}^{-2}$ for fields ~6 –16 T [22], and a significant field dependence is found for effective masses of carriers, particularly strong for β and α sheets of the Fermi surface [33].

In our study of field-tuned quantum criticality in $CeCoIn_5$ upon approaching from both inside the superconducting dome from below H_{c2} and outside the superconducting dome from above H_{c2} [46], we noticed that the heaviest electrons are strongly bound in the superconducting condensate. In this

case, the big difference in the strength of electron-electron renormalization between $CeCoIn_5$ and $CeIrIn_5$ can be directly linked to the fivefold difference in the value of the superconducting T_c. $CeCoIn_5$ and $CeIrIn_5$ have identical Fermi surfaces [32], yet interestingly they demonstrate very different electron renormalization and T_c.

Another interesting observation is that the transition at H_{c2} in $CeIrIn_5$ is second order [35], while in $CeCoIn_5$ it is first order. The first-order character of the transition is taken as arising from paramagnetic limiting of the upper critical field. We can speculate that orbital limiting of H_{c2} observed in $CeIrIn_5$ in configuration $H \parallel c$ may be responsible for the lack of a coincident quantum critical point.

Summarizing, we do not find a field-tuned quantum critical point at bulk H_{c2} in $CeIrIn_5$, nor do we see indications of its existence in all ranges of magnetic fields studied (up to 11 T). This is in sharp contrast to $CeCoIn_5$, suggesting that $CeIrIn_5$ is significantly farther away from the QCP. This finds support in the tenfold difference of the Fermi liquid T^2 coefficients between the two compounds.

C. Kadowaki-Woods ratio

It is of interest to note that the very big difference in A and B coefficients between $CeCoIn_5$ and $CeIrIn_5$ is not expected based on previous specific-heat studies [34–36]. In zero field, both materials have similar $\gamma = C/T$ values above T_c, somewhat higher in $CeInIn_5$. The big difference in A coefficients suggests that the Kadowaki-Woods ratio, defined as $r_{KW} = A/\gamma^2$, is anomalously low in $CeIrIn_5$. Indeed, as shown in Fig. 3, near H_{c2} in $CeIrIn_5$ we find $r_{KW} \sim 0.1a_0$ ($a_0 = 10\ \mu\Omega\,\mathrm{cm\,mol^2\,K^2/J^2}$), and the ratio does not show a significant field dependence. In $CeCoIn_5$, for instance, this ratio is about $0.52a_0$ at $H_{c2} \sim 6$ T, which falls on the universal line for some other heavy fermions [22].

The Kadowaki-Woods ratio cancels out mass renormalization in electrical resistivity and heat capacity measurements and leads to scaling of A and γ^2 proportional to level degeneracy of f electrons [58]. Ce ions in $CeCoIn_5$ and $CeIrIn_5$ are in the same crystal field environment, so it is natural to take the same value $n = 2$. While the r_{KW} for $CeCoIn_5$ has a value typical for heavy-fermion materials, it is anomalously low in $CeIrIn_5$.

The anomalously low value of the Kadowaki-Woods ratio in $CeIrIn_5$ suggests that different groups of carriers are determining the electrical resistivity and the heat capacity of the samples. Two scenarios can be at play. In the first, conduction electrons are strongly disconnected from f electrons due to the localized character of the latter. In this case, similar values of heat capacity in $CeCoIn_5$ and $CeIrIn_5$ reflect a nearly identical f-electron contribution, but a significantly more localized character in $CeIrIn_5$. In the second scenario, we need to assume a significant difference in the properties of different groups of carriers on the multiband Fermi surface of the materials. This scenario would require the presence of light and mobile carriers dominating in the electrical transport and heavy carriers dominating the heat capacity.

One possible way to distinguish between these two scenarios is to study magnetoresistance. In the latter scenario, the magnetoresistance of $CeIrIn_5$ should be significantly larger than in $CeCoIn_5$. Comparing the resistivity data for $CeIrIn_5$ in a magnetic field of 11 T in Fig. 1 (bottom-left panel) with similar data in $CeCoIn_5$ [22], we find that orbital magnetoresistance at the lowest temperatures is stronger in $CeCoIn_5$, which is at odds with this scenario.

V. CONCLUSIONS

In conclusion, temperature-dependent electrical resistivity and thermal conductivity measurements in $CeIrIn_5$ as functions of magnetic fields applied along the c axis revealed the validity of the Wiedemann-Franz law in the $T \to 0$ limit as long as resistivity measurements are not affected by filamentary superconductivity. At temperatures above 1 K, the temperature-dependent resistivity is well represented by a power-law function $\Delta\rho \propto T^n$, with nearly field-independent $n = 1.21$. At the lowest temperatures, the temperature dependences of both electrical resistivity, $\rho(T)$, and thermal resistivity, $\rho(T)$, follow the expectations of the Landau Fermi-liquid theory with $B \approx 2A$, as expected for Fermi surfaces without sharp features in the angular Fermi velocity distributions. The coefficient B does not exhibit a significant field dependence even on approaching $H_{c2} = 0.4$ T of the bulk superconducting state. The weak response to the magnetic field is in stark contrast with the behavior found in the closely related $CeCoIn_5$, in which the field-tuned quantum critical point coincides with H_{c2}. The value of the electron-electron mass enhancement, as judged by the A and B coefficients, is about one order of magnitude reduced in $CeIrIn_5$ as compared to $CeCoIn_5$, suggesting that the material is significantly farther away from the magnetic quantum critical point. The suppressed Kadowaki-Woods ratio in $CeIrIn_5$ compared to $CeCoIn_5$ suggests a notably more localized nature of f electrons in the compound.

ACKNOWLEDGMENTS

We thank J. Corbin for his assistance with the experiment, and R. Flint for useful discussions. L.T. acknowledges support from the Canadian Institute for Advanced Research, and funding from NSERC, FRQNT, CFI, and the Canada Research Chairs Program. Part of the work was carried out at the Brookhaven National Laboratory, which is operated for the U.S. Department of Energy by Brookhaven Science Associates (No. DEAC02-98CH10886). H.Sh. would like to acknowledge the Iran National Science Foundation (INSF) for supporting this project.

[1] C. Petrovic, R. Movshovich, M. Jaime, P. G. Pagliuso, M. F. Hundley, J. L. Sarrao, Z. Fisk, and J. D. Thompson, Europhys. Lett. **53**, 354 (2001).

[2] C. Petrovic, P. Pagliuso, M. Hundley, R. Movshovich, J. Sarrao, J. D. Thompson, Z. Fisk, and P. J. Monthoux, J. Phys.: Condens. Matter **13**, L337 (2001).

[3] J. Paglione, P.-C. Ho, M. B. Maple, M. A. Tanatar, L. Taillefer, Y. Lee, and C. Petrovic, Phys. Rev. B **77**, 100505(R) (2008).
[4] G. F. Chen, K. Matsubayashi, S. Ban, K. Deguchi, and N. K. Sato, Phys. Rev. Lett. **97**, 017005 (2006).
[5] H. Hegger, C. Petrovic, E. G. Moshopoulou, M. F. Hundley, J. L. Sarrao, Z. Fisk, and J. D. Thompson, Phys. Rev. Lett. **84**, 4986 (2000).
[6] L. D. Pham, T. Park, S. Maquilon, J. D. Thompson, and Z. Fisk, Phys. Rev. Lett. **97**, 056404 (2006).
[7] N. D. Mathur, F. M. Grosche, S. R. Julian, I. R. Walker, D. M. Freye, R. K. W. Haselwimmer, and G. G. Lonzarich, Nature (London) **394**, 39 (1998).
[8] B. D. White, J. D. Thompson, and M. B. Maple, Physica C **514**, 246 (2015).
[9] J. L. Sarrao and J. D. Thompson, J. Phys. Soc. Jpn. **76**, 051013 (2007).
[10] T. Shang, R. E. Baumbach, K. Gofryk, F. Ronning, Z. F. Weng, J. L. Zhang, X. Lu, E. D. Bauer, J. D. Thompson, and H. Q. Yuan, Phys. Rev. B **89**, 041101(R) (2014).
[11] P. G. Pagliuso, C. Petrovic, R. Movshovich, D. Hall, M. F. Hundley, J. L. Sarrao, J. D. Thompson, and Z. Fisk, Phys. Rev. B **64**, 100503 (2001).
[12] Y. Chen, W. B. Jiang, C. Y. Guo, F. Ronning, E. D. Bauer, Tuson Park, H. Q. Yuan, Z. Fisk, J. D. Thompson, and X. Lu, Phys. Rev. Lett. **114**, 146403 (2015).
[13] V. A. Sidorov, M. Nicklas, P. G. Pagliuso, J. L. Sarrao, Y. Bang, A. V. Balatsky, and J. D. Thompson, Phys. Rev. Lett. **89**, 157004 (2002).
[14] M. A. Tanatar, J. Paglione, S. Nakatsuji, D. G. Hawthorn, E. Boaknin, R. W. Hill, F. Ronning, M. Sutherland, L. Taillefer, C. Petrovic, P. C. Canfield, and Z. Fisk, Phys. Rev. Lett. **95**, 067002 (2005).
[15] H. Shakeripour, M. A. Tanatar, S. Y. Li, C. Petrovic, and L. Taillefer, Phys. Rev. Lett. **99**, 187004 (2007).
[16] H. Shakeripour, M. A. Tanatar, C. Petrovic, and L. Taillefer, Phys. Rev. B **82**, 184531 (2010).
[17] Y. Matsuda, K. Izawa, and I. Vekhter, J. Phys.: Condens. Matter **18**, R705 (2006).
[18] K. An, T. Sakakibara, R. Settai, Y. Onuki, M. Hiragi, M. Ichioka, and K. Machida, Phys. Rev. Lett. **104**, 037002 (2010); H. Aoki, T. Sakakibara, H. Shishido, R. Settai, Y. Onuki, P. Miranovic, and K. Machida, J. Phys.: Condens. Matter **16**, L13 (2004).
[19] M. P. Allan, F. Massee, D. K. Morr, J. van Dyke, A. W. Rost, A. P. Mackenzie, C. Petrovic, and J. C. Davis, Nat. Phys. **9**, 468 (2013).
[20] B. B. Zhou, S. Misra, E. H. da Silva Neto, P. Aynajian, R. E. Baumbach, J. D. Thompson, E. D. Bauer, and A. Yazdani, Nat. Phys. **9**, 474 (2013).
[21] A. Bianchi, R. Movshovich, I. Vekhter, P. G. Pagliuso, and J. L. Sarrao, Phys. Rev. Lett. **91**, 257001 (2003).
[22] J. Paglione, M. A. Tanatar, D. G. Hawthorn, E. Boaknin, R. W. Hill, F. Ronning, M. Sutherland, L. Taillefer, C. Petrovic, and P. C. Canfield, Phys. Rev. Lett. **91**, 246405 (2003).
[23] F. Ronning, C. Capan, A. Bianchi, R. Movshovich, A. Lacerda, M. F. Hundley, J. D. Thompson, P. G. Pagliuso, and J. L. Sarrao, Phys. Rev. B **71**, 104528 (2005).
[24] M. A. Tanatar, J. Paglione, C. Petrovic, and L. Taillefer, Science **316**, 1320 (2007).
[25] R. Settai, H. Shishido, S. Ikeda, Y. Murakawa, M. Nakashima, D. Aoki, Y. Haga, H. Harima, and Y. Onuki., J. Phys.: Condens. Matter **13**, L627 (2001).
[26] A. McCollam, S. R. Julian, P. M. C. Rourke, D. Aoki, and J. Flouquet, Phys. Rev. Lett. **94**, 186401 (2005).
[27] J. Paglione, M. A. Tanatar, D. G. Hawthorn, F. Ronning, R. W. Hill, M. Sutherland, L. Taillefer, and C. Petrovic, Phys. Rev. Lett. **97**, 106606 (2006).
[28] G. Knebel, D. Aoki, D. Braithwaite, B. Salce, and J. Flouquet, Phys. Rev. B **74**, 020501 (2006); G. Knebel, D. Aoki, and J. Flouquet, C. R. Phys. **12**, 542 (2011).
[29] J. K. Dong, H. Zhang, X. Qiu, B. Y. Pan, Y. F. Dai, T. Y. Guan, S. Y. Zhou, D. Gnida, D. Kaczorowski, and S. Y. Li, Phys. Rev. X **1**, 011010 (2011).
[30] N. P. Butch, K. Jin, K. Kirshenbaum, R. L. Greene, and J. Paglione, Proc. Natl. Acad. Sci. (USA) **109**, 8440 (2012).
[31] J. K. Dong, S. Y. Zhou, T. Y. Guan, H. Zhang, Y. F. Dai, X. Qiu, X. F. Wang, Y. He, X. H. Chen, and S. Y. Li, Phys. Rev. Lett. **104**, 087005 (2010).
[32] S. Elgazzar, I. Opahle, R. Hayn, and P. M. Oppeneer, Phys. Rev. B **69**, 214510 (2004).
[33] H. Shishido, R. Settai, D. Aoki, S. Ikeda, H. Nakawaki, N. Nakamura, T. Iizuka, Y. Inada, K. Sugiyama, T. Takeuchi *et al.*, J. Phys. Soc. Jpn. **71**, 162 (2002).
[34] J. S. Kim, J. Alwood, G. R. Stewart, J. L. Sarrao, and J. D. Thompson, Phys. Rev. B **64**, 134524 (2001).
[35] Y. Aoki, S. Kittaka, T. Sakakibara, A. Sakai, S. Nakatsuji, Y. Tsutsumi, M. Ichioka, and K. Machida, J. Phys. Soc. Jpn. **81**, SB014 (2012).
[36] C. Capan, A. Bianchi, F. Ronning, A. Lacerda, J. D. Thompson, M. F. Hundley, P. G. Pagliuso, J. L. Sarrao, and R. Movshovich, Phys. Rev. B **70**, 180502(R) (2004).
[37] C. Capan, L. Balicas, T. P. Murphy, E. C. Palm, R. Movshovich, D. Hall, S. W. Tozer, M. F. Hundley, E. D. Bauer, J. D. Thompson, J. L. Sarrao, J. F. DiTusa, R. G. Goodrich, and Z. Fisk, Phys. Rev. B **80**, 094518 (2009).
[38] J. S. Kim, J. Alwood, P. Kumar, and G. R. Stewart, Phys. Rev. B **65**, 174520 (2002).
[39] S. Nair, M. Nicklas, F. Steglich, J. L. Sarrao, J. D. Thompson, A. J. Schofield, and S. Wirth, Phys. Rev. B **79**, 094501 (2009).
[40] S. Nair, S. Wirth, M. Nicklas, J. L. Sarrao, J. D. Thompson, Z. Fisk, and F. Steglich, Phys. Rev. Lett. **100**, 137003 (2008).
[41] A. Bianchi, R. Movshovich, M. Jaime, J. D. Thompson, P. G. Pagliuso, and J. L. Sarrao, Phys. Rev. B **64**, 220504(R) (2001).
[42] J. Paglione, M. A. Tanatar, D. G. Hawthorn, R. W. Hill, F. Ronning, M. Sutherland, L. Taillefer, C. Petrovic, and P. C. Canfield, Phys. Rev. Lett. **94**, 216602 (2005).
[43] S. Wirth, Y. Prots, M. Wedel, S. Ernst, S. Kirchner, Z. Fisk, J. D. Thompson, F. Steglich, and Y. Grin, J. Phys. Soc. Jpn. **83**, 061009 (2014).
[44] Y. Maeno, T. Ando, Y. Mori, E. Ohmichi, S. Ikeda, S. NishiZaki, and S. Nakatsuji, Phys. Rev. Lett. **81**, 3765 (1998).
[45] R. Hu, K. Lauritch-Kullas, J. O'Brian, V. F. Mitrovic, and C. Petrovic, Phys. Rev. B **75**, 064517 (2007).
[46] J. Paglione, M. A. Tanatar, J.-Ph. Reid, H. Shakeripour, C. Petrovic, and L. Taillefer, arXiv:1406.0031.
[47] J. A. Hertz, Phys. Rev. B **14**, 1165 (1976); A. J. Millis, *ibid.* **48**, 7183 (1993).
[48] A. Rosch, Phys. Rev. B **62**, 4945 (2000).

[49] T. Moriya, Y. Takahashi, and K. Ueda, J. Phys. Soc. Jpn. **59**, 2905 (1990).

[50] P. Gegenwart, J. Custers, C. Geibel, K. Neumaier, T. Tayama, K. Tenya, O. Trovarelli, and F. Steglich, Phys. Rev. Lett. **89**, 056402 (2002).

[51] J. Custers, P. Genewart, H. Wilhelm, K. Neumaier, Y. Tokiwa, O. Trovarelli, C. Geibel, F. Steglich, C. Pepin, and P. Coleman, Nature (London) **424**, 524 (2003).

[52] K. Heuser, J. S. Kim, E. W. Scheidt, T. Schreiner, and G. R. Stewart, Physica B **261**, 392 (1999).

[53] A. Malinowski, M. F. Hundley, C. Capan, F. Ronning, R. Movshovich, N. O. Moreno, J. L. Sarrao, and J. D. Thompson, Phys. Rev. B **72**, 184506 (2005).

[54] D. K. Wagner, J. C. Garland, and R. Bowers, Phys. Rev. B **3**, 3141 (1971).

[55] G. K. White and R. J. Tainsh, Phys. Rev. Lett. **19**, 165 (1967).

[56] B. Lussier, B. Ellman, and L. Taillefer, Phys. Rev. Lett. **73**, 3294 (1994).

[57] Y. Haga, Y. Inada, H. Harima, K. Oikawa, M. Murakawa, H. Nakawaki, Y. Tokiwa, D. Aoki, H. Shishido, S. Ikeda, N. Watanabe, and Y. Onuki, Phys. Rev. B **63**, 060503 (2001).

[58] P. Coleman, *Heavy Fermions: Electrons at the Edge of Magnetism, Handbook of Magnetism and Advanced Magnetic Materials*, edited by H. Kronmuller and S. Parkin, Vol. 1: Fundamentals and Theory (John Wiley and Sons, 2007), pp. 95–148.