

Structure and physical properties of SrNiRu₅O₁₁ single crystals: An *R*-type ferrite based on ordered kagome nets

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Single crystals of the *R*-type ferrite SrNiRu₅O₁₁ were grown from a chloride flux. The hexagonal crystal structure contains ruthenium located on distorted kagome nets. The low-temperature dc magnetic susceptibilities (χ_{\perp} and χ_{\parallel} , perpendicular and parallel to the *c* axis, respectively) diverge as $T^{-0.3}$, and do not exhibit any indication of long-range magnetic order down to 4.5 K. The electrical resistivity varies as $T^{1.6}$ below 40 K, which is typical of non-Fermi liquids, and may originate from a competition between residual magnetic interactions among Ni²⁺ ($S = 1$) spins and geometrical frustration on the two-dimensional kagome lattice of Ru³⁺ ($S = 1/2$) spins. The transverse magnetoresistivity ρ_{xy} at constant temperature $T = 5$ K for current (*J*) -magnetic field (*H*) configurations, $\mathbf{J} \perp \mathbf{H} \parallel \mathbf{c}$ axis and $\mathbf{J} \parallel \mathbf{H} \perp \mathbf{c}$ axis, reveals no anomalous contribution, which is consistent with the absence of magnetic order. Fits of the specific heat data below 10 K require a dominant, but unusual electronic term of the form $C_{el} = \gamma T^{1.2}$, which is expected for massless Dirac fermion states in topological insulators, or spin-liquid phases.

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I. INTRODUCTION

Ruthenium oxides have been intensively investigated over the last two decades as a very interesting class of transition metal oxides with unusual physical properties [1]. One of the most important features of 4*d* transition metal oxides is the large spatial extent of their 4*d*-electron orbitals, which favors their greater delocalization compared to their 3*d* counterparts, which leads to increased hybridization of transition metal and oxygen states, and modifies crystalline electric-field interactions. These conditions are known to promote a variety of exotic magnetic properties in ruthenium oxides [2].

Previously, we have studied ternary ruthenium ferrites with chemical formula (Ba,Sr)T_{2±x}Ru_{4±x}O₁₁ (T = Fe, Co, Mn), which crystallize in the hexagonal *R*-type ferrite structure [3–7], which contains a kagome substructure of mixed occupancy by Ru and transition element T, as shown in Fig. 1. Layers of edge-sharing M(2)O₆ octahedra form a kagome net within the *ab* plane; and the kagome planes are connected along [001] by face-sharing M(1)O₆ octahedra and M(3)O₅ trigonal bipyramids. It is important to note that the variable occupation of octahedrally coordinated M(1) and M(2) sites by 3*d* elements and 4*d* ruthenium generate large homogeneity ranges, while the M(3) sites with trigonal-pyramidal coordination are exclusively occupied by 3*d* elements. The substitution of 3*d* elements for 4*d* ruthenium within large homogeneity ranges permits precise control of magnetic order (including attainment of unusually high Curie temperatures $T_C \sim 490$ K), a narrow semiconducting gap ($\Delta \sim 0.1$ eV), and a sizable anomalous Hall effect, making these materials very attractive for spintronic applications [4]. Furthermore, the mixed occupancy of the M(2) sites by Ru and M in the frustrated kagome plane gives rise to exotic physical phenomena such as nonzero scalar spin chirality [6], which has attracted increasing attention due to its potential exploitation in applications.

In this paper we report electric, magnetic, and thermodynamic properties of SrNiRu₅O₁₁ single crystals, which is a new representative of the broad family of ternary ruthenium ferrites. In contrast to most previously studied members of the *R*-type ferrites, no 3*d*/4*d* atomic disorder is observed on the M(2)/M(3) sites of SrNiRu₅O₁₁. However, the kagome net is distorted due to the formation of Ru-Ru dumbbells via metal-metal bonding in SrNiRu₅O₁₁ and BaZnRu₅O₁₁ [8]. Together with the compounds (Ba,Sr)TRu₅O₁₁ with T = Li, Zn, and Cu [8–10], SrNiRu₅O₁₁ constitutes an uncommon example of an ordered structure within the *R*-type ferrite family. Our results verify that such an ordered structure enhances magnetic frustration within the kagome planes, and allows one to study a proximate quantum critical state that generates non-Fermi-liquid (NFL) behavior at low temperatures (also observed for BaZnRu₅O₁₁ [9]).

II. EXPERIMENTAL DETAILS

The dc magnetic moment *m* of oriented single crystals was measured over a temperature range $5 \text{ K} \leq T \leq 300 \text{ K}$ in applied magnetic fields $0 \leq \mu_0 H \leq 5 \text{ T}$ using a Quantum Design MPMS5 SQUID magnetometer. Longitudinal and transverse magnetoresistivities, $\rho_{xx}(T)$ and $\rho_{xy}(H, T)$, respectively, were measured using a dc four-probe setup with currents *J* directed in the *ab* plane with magnitudes $5 \text{ mA} \leq J \leq 20 \text{ mA}$. The specific heats of several crystals were measured over a temperature range 0.35–50 K using a Quantum Design PPMS9 system.

III. RESULTS AND DISCUSSION

A. Crystal structure and composition

We have previously reported on the growth of SrNiRu₅O₁₁ single crystals (Fig. 2) from SrCl₂ flux, and their crystal structure determination [8]. According to microprobe chemical analysis and x-ray structure refinements, the composition is well represented by SrNiRu₅O₁₁ with no indication of any significant homogeneity range. However, SrNiRu₅O₁₁

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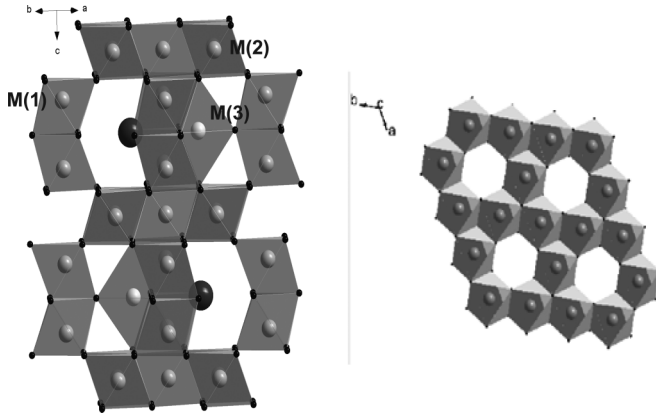


FIG. 1. Left panel: Crystal structure of $\text{SrNiRu}_5\text{O}_{11}$. Dark dots are oxygen atoms and large gray spheres are transition elements (Ru or Ni). The M(1) ($4e$) sites in face-shared octahedra, as well as M(2) ($6g$) sites in edge-shared octahedra in the kagome sublattice, are both occupied by Ru. The M(3) ($2d$) sites in double trigonal prisms are exclusively occupied by Ni. Dark spheres are Sr sites. Right panel: Arrangement of edge-sharing coordination polyhedra of transition metal ions in the ab (kagome) plane.

exhibits a doubled unit cell parameter a [$a = 11.7374(2)$ Å, $c = 13.2088(4)$ Å] with a lower space group symmetry, $P6_3/m$, as compared to the $P6_3/mmc$ R -type hexagonal unit cell [8]. Figure 1 shows a section of the smaller R -type substructure, which is sufficient for discussion of the metal atom sites. Ni atoms are restricted to the trigonal bipyramids of oxide ions, but randomly occupy either one or the other of the two adjacent tetrahedrally coordinated sites within the bipyramids. The Ru(2) sites have octahedral coordination by oxygen, and form distorted kagome nets, whereas the Ru(1) occupy face-sharing, double octahedra that connect the kagome nets along the $[001]$ direction.

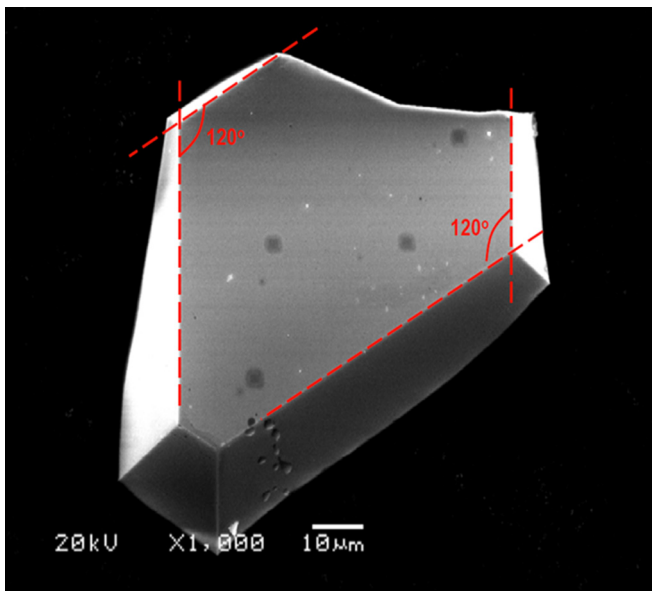


FIG. 2. Scanning electron microscopy image of a small crystal of $\text{SrNiRu}_5\text{O}_{11}$, revealing a hexagonal crystal habit, as shown by the characteristic angles highlighted by red dashed lines.

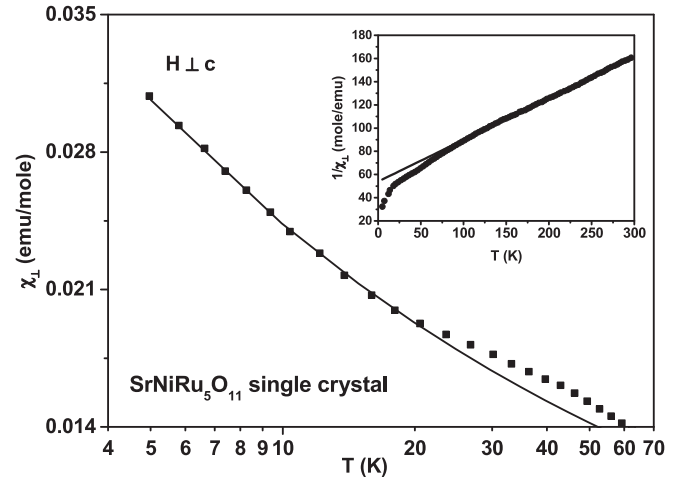


FIG. 3. Low-temperature magnetic susceptibility for field $\mu_0 H = 0.5$ T applied perpendicular to the c axis, χ_{\perp} , versus temperature T . The solid curve is a fit to $\chi_{\perp} \sim T^{-0.3}$. Inset: The inverse magnetic susceptibility $(\chi_{\perp})^{-1}$ versus T ; the solid curve is a fit of the data to the Curie-Weiss law (see text for details).

B. Magnetic and transport properties

The magnetic species in $\text{SrNiRu}_5\text{O}_{11}$ include Ru^{3+} [$\text{Ru}(2), S = 1/2$] ions that form a kagome sublattice in the ab planes, which are separated from adjacent kagome planes by magnetic Ni^{2+} ($S = 1$) and nonmagnetic Ru^{4+} [$\text{Ru}(1), S = 0$] ions. The proposed spin configurations for Ru and Ni are based on the crystal structure refinements, which suggest a $3d^8$ electron configuration for the Ni [M(3) site], $4d^5$ for the Ru(2) [M(2) site], and $4d^4$ for the Ru(1) [M(1) site] [9]. We assume Ru(1) to be in a low-spin state ($\text{Ru}^{4+}, S = 0$) since neutron diffraction investigations performed on Co- and Mn-bearing ferrites consistently indicate there is no magnetic moment on this site [10,11].

Our high-temperature magnetic susceptibility data are shown in the insets to Figs. 3 and 4. The inverse field-cooled dc susceptibility for $H \perp c$ (Fig. 3, inset), $1/\chi_{\perp}$, presents a change in slope near 60 K and a near-linear temperature dependence above 100 K. A Curie-Weiss fit of the data in the temperature interval $100 \text{ K} < T < 300 \text{ K}$ gives a Curie-Weiss temperature $\theta_p = -150 \text{ K}$ and an effective magnetic moment $\mu_{\text{eff}} \approx 4.7 \mu_B$. The susceptibility of single-crystal $\text{SrNiRu}_5\text{O}_{11}$ is anisotropic and about two times higher for $H \parallel c$ than for $H \perp c$. Above 100 K, χ_{\parallel} follows a Curie-Weiss law with $\theta_p = -69 \text{ K}$, and $\mu_{\text{eff}} = 4.9 \mu_B$ (Fig. 4, inset). The negative θ_p value is consistent with the dominance of antiferromagnetic interactions between room temperature and approximately 100 K. An estimation of the theoretical effective magnetic moment can be obtained using an expression $\mu_{\text{eff}} = [n_1 \mu^2(\text{Ni}^{2+}) + n_2 \mu^2(\text{Ru}^{3+}) + n_3 \mu^2(\text{Ru}^{4+})]^{1/2}$, where n_1 , n_2 , and n_3 are the molar fractions of Ni^{2+} , Ru^{3+} , and Ru^{4+} , respectively. Insertion of the spin-only effective magnetic moments μ for Ru^{3+} ($S = 1/2$), Ru^{4+} ($S = 0$), and Ni^{2+} ($S = 1$) yields a value for the net effective moment, $\mu_{\text{eff}} = 4.1 \mu_B$, which is in reasonable agreement with the experimental values. On the other hand, the experimental value of μ_{eff} is close to the range of the values typically observed for Ni^{2+} in high-spin

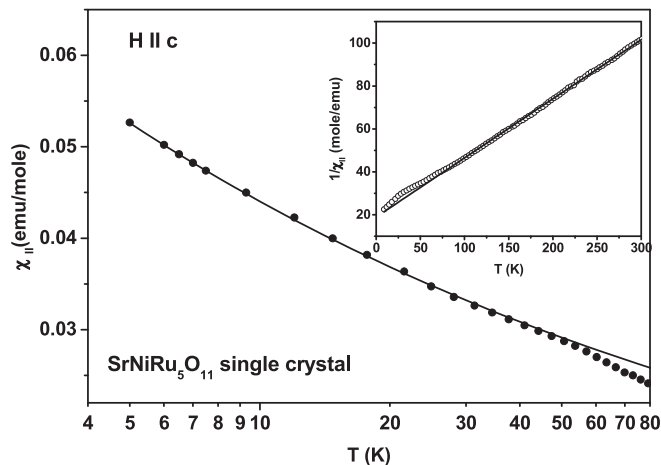


FIG. 4. Low-temperature magnetic susceptibility for magnetic field $\mu_0 H = 0.5$ T applied parallel to the c axis, χ_{\parallel} , versus temperature T . The solid curve is a fit to $\chi_{\parallel} \sim T^{-0.3}$. Inset: The inverse magnetic susceptibility $(\chi_{\parallel})^{-1}$ versus T ; the solid curve is a fit of the data to the Curie-Weiss law (see text for details).

complexes [12], and suggests the presence of an orbital contribution. For example, in tetrahedral coordination, Ni^{2+} favors the ${}^3T_{1g}$ triplet state, which may be subject to spin-orbit coupling; these circumstances could lead to potential effective moments for Ni^{2+} lying in the range $3.5 - 4.2 \mu_B$ [12]. It must be stressed that these calculated values of the effective magnetic moment are based on rough estimates that do not take into consideration strong spin-orbit interactions present in $4d$ transition metal oxides [2].

The most important conclusion we draw from the low-temperature data for the magnetic susceptibility of single-crystal $\text{SrNiRu}_5\text{O}_{11}$ under a magnetic field of $\mu_0 H = 0.5$ T applied either parallel or perpendicular to the c axis (see Figs. 3 and 4, respectively) is that $\chi(T) \sim T^{-0.3}$ over a *one-order-of-magnitude* range of temperature. More precisely, the magnetic susceptibility follows a $T^{-0.33}$ dependence below 20 K for $H \perp c$, and a $T^{-0.26}$ dependence below 40 K for $H \parallel c$. This behavior of $\text{SrNiRu}_5\text{O}_{11}$ conflicts with simple Fermi-liquid theory, where χ tends to a constant as $T \rightarrow 0$. The low-temperature divergence of the magnetic susceptibility suggests the existence of unusual magnetic excitations that differ from those of familiar Fermi liquids. On the other hand, a $T^{-0.3}$ power law can be attributed to NFL behavior that is associated with a nearby low-temperature quantum critical point (QCP) between paramagnetic and magnetically ordered states [13]. The magnetization curves are nonlinear at 5 K for both directions of the applied field, as shown in Fig. 5. A small hysteresis is observed when field is applied parallel to the c axis, suggesting weak ferromagnetic correlations. Above 130 K, the magnetic hysteresis loops are linear (Fig. 6). Note that the neutron diffraction investigations on metallic Mn and Co analogs indicate that M(3) sites always order with spins ferromagnetically aligned along the c axis [10,11]. This motivates our speculation that Ni^{2+} ions occupying M(3) sites are responsible for weak ferromagnetic correlations evident in the sample data below 130 K.

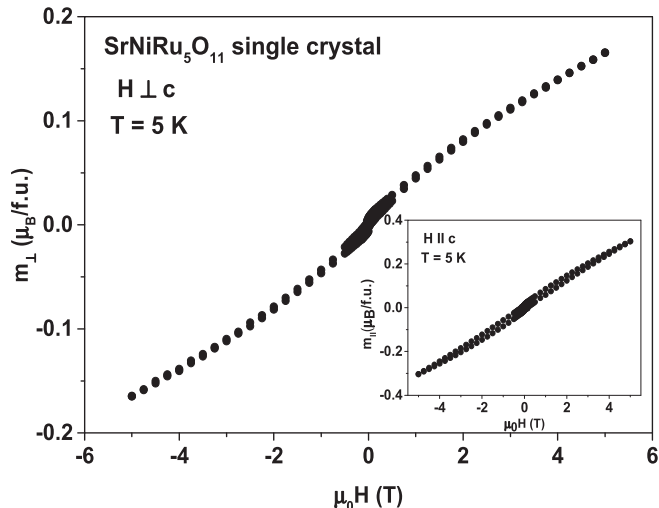


FIG. 5. Magnetic moment m_{\perp} perpendicular to the c axis versus magnetic field H at temperature $T = 5$ K. Inset: The magnetic moment m_{\parallel} along the c axis versus magnetic field at $T = 5$ K.

Our magnetic data for $\text{SrNiRu}_5\text{O}_{11}$ indicate this compound does not undergo long-range magnetic order down to 4.5 K, in contrast to the strong ferromagnetic order previously observed for $(\text{Ba,Sr})\text{T}_{2\pm x}\text{Ru}_{4\pm x}\text{O}_{11}$ ($T = \text{Fe, Co, Mn}$) [10,11]. This result points to the presence of much stronger geometrical frustration in the Ni-bearing ferrite due to the reduced disorder on Ru(2) sites in the kagome planes.

Measurements of the Hall coefficient are a good alternative probe of the magnetization of ferromagnetic materials, since the Hall resistivity is expressed as

$$\rho_{xy} = R_0 H + 4\pi M R_s, \quad (1)$$

where R_0 is the Hall coefficient resulting from the Lorentz force on the carriers, and R_s is the coefficient describing the anomalous Hall effect that depends on the magnetization

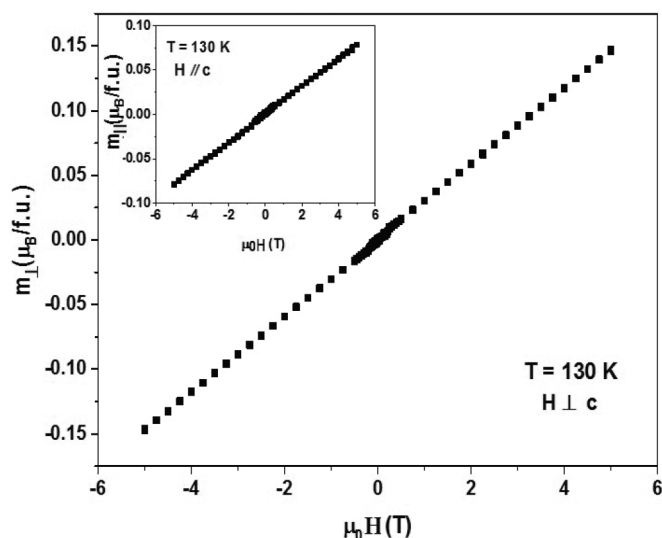


FIG. 6. Magnetic moment for applied field H perpendicular to the c axis, m_{\perp} , versus H at temperature $T = 130$ K. Inset: The moment for H perpendicular to the c axis, m_{\parallel} , versus H at $T = 130$ K.

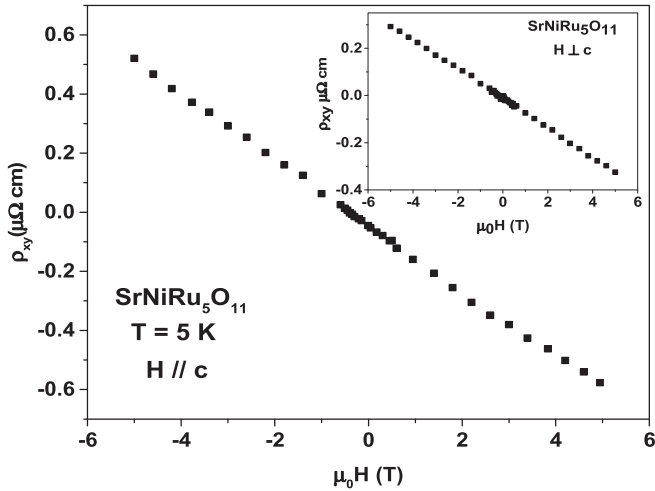


FIG. 7. Transverse resistivity ρ_{xy} versus magnetic field $H \parallel c$ axis for single-crystal $\text{SrNiRu}_5\text{O}_{11}$ at temperature $T = 5$ K. Inset: The transverse resistivity ρ_{xy} versus magnetic field $H \perp c$ axis for single-crystal $\text{SrNiRu}_5\text{O}_{11}$ at $T = 5$ K.

M and spin-orbit coupling [14,15]. In particular, ρ_{xy} has roughly the same field dependence as the magnetization in the ferromagnetic state below T_C . Moreover, a new “topological Hall effect” (THE) related to the spin chirality in noncoplanar spin configurations has been recently proposed [16]. In contrast with the conventional mechanism, the THE does not require any spin-orbit coupling, and arises solely from the Berry phase acquired by an electron moving in a smoothly varying magnetization [17]. A THE driven by spin chirality has been experimentally observed in ferromagnetic Co- and Mn-bearing ferrites [6,11]. Ferromagnetic materials are not ideal for inducing a large THE, since strong ferromagnetic correlations tend to mask the THE when field is applied along the net magnetization of the sample. Although $\text{SrNiRu}_5\text{O}_{11}$ appears to be a paramagnetic semimetal, it could alternatively have an “all-in/all-out” or “ $q = 0$ ” planar spin configuration [10] with no net moment, in which case it might exhibit a THE in an external field that induces a noncoplanar (e.g., canted along the c axis) spin configuration among the Ru ions residing on the kagome sublattice [6].

We measured $\rho_{xy}(H, T = 5 \text{ K})$ in single-crystal $\text{SrNiRu}_5\text{O}_{11}$ with an in-plane current J and two directions of the applied magnetic field, $J \perp H \parallel c$ and $J \parallel H \perp c$ at $T = 5$ K (Fig. 7). In contrast to the Co, Mn, or Fe-analogs [6,11,18], we found no anomalous Hall signal in $\text{SrNiRu}_5\text{O}_{11}$, in that ρ_{xy} varies linearly on magnetic field applied in both directions. The negative slope of $\rho_{xy}(H)$ indicates that the dominant charge carriers in $\text{SrNiRu}_5\text{O}_{11}$ are electrons with a carrier concentration $n \approx 10^{21} \text{ cm}^{-3}$ and mobility $\mu = R_0/\rho \approx 4.7 \times 10^{-2} \text{ m}^2 \text{ V}^{-1} \text{ s}^{-1}$ at $T = 5$ K [using Eq. (1) with $R_0 = d\rho_{xy}/dH = 1/nec$, where e is the electron charge and c is the speed of light). Note that the carrier concentration is one order of magnitude higher for $\text{SrNiRu}_5\text{O}_{11}$ than for Mn- or Co-bearing ferrites [6,11], indicating the density of states $N(E_F)$ at the Fermi level increases when Co or Mn are replaced by Ni. The lack of an anomalous contribution to the Hall resistivity is

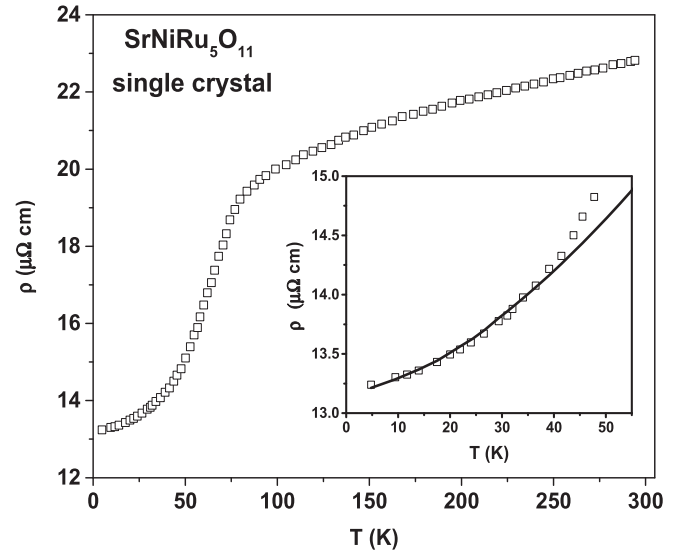


FIG. 8. Temperature dependence of the longitudinal electrical resistivity $\rho = \rho_{xx}$ of single-crystal $\text{SrNiRu}_5\text{O}_{11}$ with in-plane current $J \perp c$ axis. Inset: The low-temperature resistivity, where the solid curve is a fit of the data to $\rho_{xx} \sim T^{1.6}$.

highly significant, and suggests that any magnetic coupling developing below 130 K is probably limited to short-range correlations. Furthermore, the absence of a nonmonotonic dependence of ρ_{xy} (which is linked to a possible THE [6]) and previous magnetic neutron scattering results on Mn- and Co-bearing ferrites [10,11], argue against a canted variant of an in-plane, coplanar “ $q = 0$ ” arrangement of Ru^{3+} ($S = 1/2$) spins residing on the kagome triangles in $\text{SrNiRu}_5\text{O}_{11}$.

The temperature dependence of the electrical resistivity of single-crystal $\text{SrNiRu}_5\text{O}_{11}$ with in-plane current $J \perp H \parallel c$ is shown in Fig. 8. The compound exhibits metallic conductivity down to 4.5 K. The residual resistivity $\rho_0(T = 4.5 \text{ K}) = 13.2 \mu\Omega \text{ cm}$ indicates a reasonably good-quality sample. No significant difference is observed between measurements in zero field or 5 T. It is interesting that a near-linear temperature dependence of the in-plane resistivity persists over the high-temperature interval, 200–300 K. The strong convexity of $\rho(T)$ near 100 K is similar to the “saturating resistivities” of strongly paramagnetic actinide compounds [19]. This behavior also mimics high-temperature cuprate or iron-arsenide-based superconductors, where it is considered a signature of marginal Fermi-liquid physics. As the temperature decreases below about 200 K, the resistivity gradually deviates from linearity, and exhibits a clear slope change near 70 K. A rapid decrease of the resistivity below 70 K suggests a crossover between temperature-dependent scattering mechanisms of different types. Below 40 K, the resistivity can be fitted to a $T^{1.6}$ dependence, as shown in the inset of Fig. 8; this power law is a signature of scattering by strong spin fluctuations in metallic actinide and transition element materials that are nearly magnetic or weakly ferromagnetic [19,20]. In contrast, a T^2 dependence of the low-temperature electrical resistivity is predicted by Fermi-liquid theory in the limit $T \rightarrow 0$ K. The reduction of the resistivity exponent below 2 is expected for an antiferromagnetic quantum critical point,

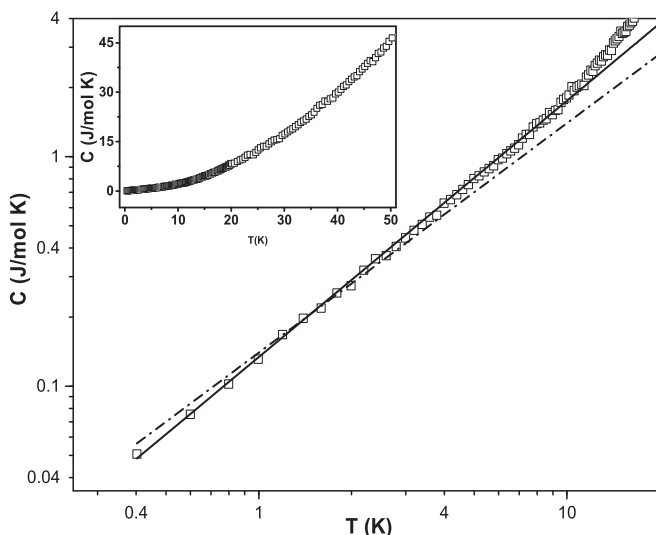


FIG. 9. Low-temperature specific heat $C(T)$ of single-crystal $\text{SrNiRu}_5\text{O}_{11}$ versus temperature T (note logarithmic scale). The dashed line is a fit of the data to Eq. (2) with power-law exponent $\alpha = 1.0$. The solid line is an improved fit using Eq. (2) with a modified electronic contribution, $C_{el} = \gamma T^{1.2}$. The inset shows the total specific heat $C(T)$ data for single-crystal $\text{SrNiRu}_5\text{O}_{11}$ over the temperature range 0.35–50 K.

since the strong scattering of electrons takes place at particular antiferromagnetic wave vectors (“hot spots”) around the Fermi surface, allowing larger areas of the Fermi surface to “short out” this strong scattering with conventional Fermi-liquid transport [21]. The deviation of the resistivity data from Fermi-liquid behavior below 40 K is therefore consistent with NFL behavior, and also consistent with the temperature dependence of the magnetic susceptibility, which crosses over from a high-temperature Curie-Weiss behavior to a divergent power-law behavior predicted by NFL models. In any case, it is clear that the transport and magnetic properties of $\text{SrNiRu}_5\text{O}_{11}$ cannot be explained within Fermi-liquid theory.

C. Specific heat

The inset in Fig. 9 shows the temperature dependence of the specific heat of $\text{SrNiRu}_5\text{O}_{11}$, which reveals no anomalous behavior that could indicate phase transitions over the temperature range 0.35–50 K. The specific heat below 10 K has been analyzed by considering the electronic, lattice, and nuclear Schottky contributions. The nuclear specific heat in $\text{SrNiRu}_5\text{O}_{11}$ can be attributed to Ru and Ni isotopes with nonzero quadrupole moment (i.e., having nuclear spin $I > 1/2$) residing at sites with axial symmetry and a nonzero electric-field gradient. In addition, unpaired d electrons of Ru and Ni can generate magnetic hyperfine fields if the atomic spin relaxation is slower than the Larmor frequency of the nuclear spin [22]. We considered isotopes (with natural abundance 1.2% for ^{61}Ni , 12.76% for ^{99}Ru , and 17.06% for ^{101}Ru) that appreciably contribute to a nuclear Schottky effect arising from quadrupole and hyperfine interactions. We fitted our data below 10 K to the expression

$$C(T) = \gamma T^\alpha + \beta T^3 + D/T^2, \quad (2)$$

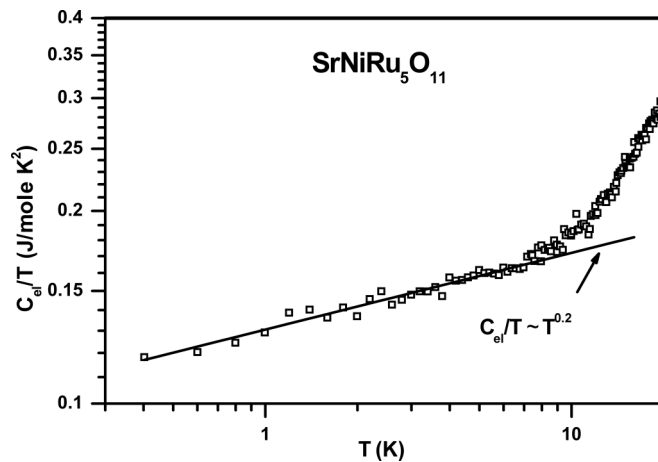


FIG. 10. Calculated electronic contribution to the specific heat C_{el} derived from the best fit of the total specific heat versus temperature, determined from the following procedure: First we obtain the best-fit description of the total specific heat data using $C(T) = \gamma T^{1.2} + \beta T^3 + D/T^2$, where the form of the first term is attributed to spin-liquid behavior [35,36]. Second, the electronic contribution was separated as a difference between the total heat capacity and the sum of nuclear Schottky and lattice contributions: $C_{el}(T) = C_{\text{total}}(T) - (\beta T^3 + D/T^2)$. We show not C_{el} , but C_{el}/T , in order to emphasize how the coefficient $\gamma \equiv C_{el}/T$ from a traditional Sommerfeld fit must decrease at the lowest temperatures, which differs from the constant value observed for conventional metals. The solid line is shown to demonstrate that C_{el}/T follows a $T^{0.2}$ law below 10 K. Note that the data are shown as a log-log plot.

where $C_{el} = \gamma T^\alpha$ is the electronic contribution, $C_{lat} = \beta T^3$ is the lattice contribution, and $C_{sch} = D/T^2$ represents the high-temperature form of the nuclear Schottky term. The standard fit using Eq. (2) with $\alpha = 1$ was relatively poor, but was substantially improved using a phenomenological C_{el} term with exponent $\alpha = 1.2$ (Fig. 9). The best-fit coefficients are $\gamma = 0.134 \text{ J/mol K}^{2.2}$, $\beta = 3.2 \times 10^{-4} \text{ J/mol K}^4$, and $D = 8.6 \times 10^{-4} \text{ J K/mol}$. Note that a fit that includes only conventional Sommerfeld electronic and Debye phonon contributions yields similar coefficients, $\gamma = 0.144 \text{ J/mol K}^2$ and $\beta = 6.5 \times 10^{-4} \text{ J/mol K}^4$.

Both types of fit indicate a relatively large electronic contribution to the specific heat for a semimetal, but typical of moderately heavy fermion materials [23]. A possible heavy fermion state in $\text{SrNiRu}_5\text{O}_{11}$ would traditionally be attributed to hybridization between localized Ni-3d and Ru-4d conduction electron states. More generally, a large γ may indicate a substantial contribution from fluctuating magnetic moments in the system. The paramagnon model [24–26] of spin fluctuations in exchange-enhanced paramagnets and nearly ferromagnetic materials predicts a heat capacity contribution of the form $T^3 \ln(T/T_{sf})$ to the electronic specific heat, where T_{sf} is a characteristic spin fluctuation temperature. However, all our attempts to include this term failed; and we therefore concluded that the fractional power-law term most adequately describes the data. The calculated electronic contribution to the total specific heat divided by T versus temperature, according to the corrected Eq. (2) that includes $C_{el} = \gamma T^{1.2}$, is shown in Fig. 10, where it is clear the behavior of C/T for $T < 10$ K

is not constant as in normal metals, but varies as $T^{0.2}$. In conventional metals $C_{el} \sim \gamma T$ at low temperatures and the clear deviation (taking into account a nuclear Schottky term to the total heat capacity) of the exponent α from unity is another indication of the unconventional character of the low-energy excitations in $\text{SrNiRu}_5\text{O}_{11}$.

Nevertheless, the unconventional power-law dependence of $C_{el}(T)$ is consistent with that expected for the massless Dirac fermion states in topological insulators (TIs), or gapless spin liquids realized on $S = 1/2$ kagome nets [27,28]. In contrast, the low-temperature electronic specific heat of $\text{BaCo}_2\text{Ru}_4\text{O}_{11}$, which is one of the lower- T_C magnetic phases of the R -type ferrite class, exhibits conventional $C_{el} \sim \gamma T$ behavior with a moderate $\gamma \approx 120 \text{ mJ/mol K}^2$ [4,29]. The observed evolution of the specific heat as a function of the $3d$ element suggests a transition from a “trivial” to “topological” behavior within the structural family of the ruthenium ferrites, which is driven by the $3d/4d$ -element ratio—in other words, by the relative strength of spin-orbit ($4d$ -element) and electron-electron ($3d$ -element) interactions. The protection of the Dirac point by time-reversal symmetry can be lifted by magnetic ordering, which takes place when Ni is substituted by Fe, Mn, or Co. As a result, a gap that separates the upper and lower branches of the Dirac cone is formed, leading to massive fermion states [30]. Note that a kagome lattice has been argued to play an essential role in the formation of a two-dimensional TI [31]. However, these scenarios are complicated by the presence of substantial atomic disorder on the kagome lattice in the magnetically ordered R -type ferrite compositions studied to date.

Based on these results one may suggest that an intermediate topological Mott insulator state exists between the topological band insulators and the familiar Mott insulating phases for relatively strong electron-electron interactions for which magnetic ordering is expected. A model based on the slave-rotor approximation and the strong coupling limit [27] yields a phase diagram with four main phases whose stabilities depend upon the Hubbard repulsion U and spin-orbit coupling λ for moderately strong electron-electron repulsion: (1) metallic, (2) topological band insulator, (3) topological Mott insulator, and (4) gapless Mott insulator. The experimental data suggest that $\text{SrNiRu}_5\text{O}_{11}$ might be located somewhere near the boundary between the gapless Mott insulator and metallic states. It should be noted that in bulk TI samples [32–34], the transport properties of the surface states are often mixed with those of bulk states, making it difficult to observe insulating behavior in $\rho(T)$. The bulk conductivity decreases while the conductivity due to the surface states increases with decreasing temperature, and the electrical conduction at low temperature is dominated by the surface. The $\rho(T)$ data of $\text{SrNiRu}_5\text{O}_{11}$ therefore could be interpreted to reflect metallic behavior that is similar to other bulk TI.

Alternatively, the enhanced Sommerfeld coefficient of $\text{SrNiRu}_5\text{O}_{11}$ and a metallic/semimetallic temperature dependence of the electrical resistivity are consistent with incipient band magnetism in this compound, suggesting that this material is an example of a frustrated system with a spin-liquid state that involves itinerant electrons. On the other hand, our observations of high values of the effective magnetic moment classify $\text{SrNiRu}_5\text{O}_{11}$ as a localized moment system. The relatively high Weiss temperatures (-150 K and -69 K

for the two crystallographic orientations) and the absence of any magnetic order down to 0.4 K may also signal a spin-liquid state, as might be expected for the kagome net of Ru^{3+} ($S = 1/2$) spins. It is interesting to note that the $C_{el} = \gamma T^\alpha$ term with exponent $\alpha = 1.2 - 1.3$ is observed in gapless spin liquids (e.g., herbertsmithite [35] and vanadium oxyfluoride [36]), where spin fluctuations persist down to $T = 0 \text{ K}$ due to the strong geometrical frustration present on the $S = 1/2$ kagome net. These insulating materials with highly localized magnetic moments nevertheless exhibit a near-linear specific heat that normally indicates the low-temperature thermal excitations are dominated by itinerant electrons. On the other hand, the physics of magnetic frustration in metallic systems is relatively unexplored, mainly due to the lack of model systems. A possible spin-liquid state has been detected in a few metallic materials that do not have a kagome net—e.g., $\text{Y}(\text{Sc})\text{Mn}_2$ with itinerant $3d$ -electron spins on a pyrochlore lattice [37]. Moreover, a number of frustrated metallic materials in a spin-liquid state exhibit a conventional linear specific heat $C \sim T$ [38].

Generally speaking, the nonanalytic behavior in both heat capacity and magnetic susceptibility, together with a $T^{1.6}$ dependence of the electrical resistivity data, indicate that the system is on the verge of a magnetic phase transition that must occur below 0.35 K (our low-temperature limit in heat capacity experiments) in the case of $\text{SrNiRu}_5\text{O}_{11}$. Given the Ru moments lie on a frustrated kagome lattice, magnetic order is likely suppressed by quantum fluctuations above a highly degenerate quantum-spin-liquid ground state, which is one type of NFL state widely discussed in literature.

In the case of intermetallic compounds, current NFL theories can be divided into three general categories: multichannel Kondo models [39], systems proximate to QCP treated by renormalization group theory of spin fluctuations [40], and models based on atomic disorder [41]. Our magnetic susceptibility and electrical resistivity data for $\text{SrNiRu}_5\text{O}_{11}$ show that these NFL models do not adequately describe our data, for the following reasons: (1) three-dimensional renormalization group theory for well-ordered materials near a QCP predicts $\chi \sim T^{3/2}$, $\rho \sim T^{3/2}$ (antiferromagnetic QCP), or $\rho \sim T^{5/3}$ (ferromagnetic QCP); (2) the multichannel Kondo or Kondo disorder models demand $\chi \sim -\ln T$ and $\rho \sim 1 - aT$; and (3) models of Kondo alloys require atomic disorder and large residual resistivities [41]. Furthermore, Kondo and atomic disorder models do not provide self-consistent agreement with our resistivity and susceptibility data, due to the fact that the residual ($T = 0$) resistivity of our single crystal is $13.2 \mu\Omega \text{ cm}$ (which indicates a well-ordered, clean sample). Moreover, $\text{SrNiRu}_5\text{O}_{11}$ has no significant homogeneity range and is structurally ordered, according to our single-crystal structure refinements [8]. We conclude that an alternative NFL mechanism is needed to explain the observed low-temperature behavior of the magnetic, transport, and thermodynamic properties of $\text{SrNiRu}_5\text{O}_{11}$.

Note that the above-cited NFL theories have been developed for f -electron, heavy fermion intermetallics, the magnetic and electronic properties of which are dictated by interactions between localized f orbitals with strong magnetic moments and itinerant electrons. In some cases, transition metal oxides also behave as heavy fermion metals [42–44] if a paramagnetic state survives at low temperatures due to geometrical frustration

[45]. Furthermore, in metallic systems magnetic order can be suppressed by frustrated interactions or by conduction electron screening through the Kondo effect. The *combination* of strong geometrical frustration and Kondo screening leads to a phase transition into a quantum-spin-liquid state that persists to very low temperatures (see [38], and references therein).

It is instructive to view SrNiRu₅O₁₁ as an ordered example of the *R*-type ferrites with strong magnetic frustration generated within the *ab* (kagome) nets. In contrast, the isostructural, metallic Co ($T_C = 105$ K) and Mn ($T_C = 183$ K) ferrites suffer from substantial *3d/4d* disorder on M(2) sites within the kagome nets [6,11]. Consequently, the magnetic frustration is depressed, and long-range ferromagnetic ordering occurs at remarkably high temperatures [4]. In the case of SrNiRu₅O₁₁, the structural disorder in the kagome planes is absent. Furthermore, in ferrites where Co and Mn are completely replaced by Ni and Ru, the electrons become more itinerant, most probably due to an increased density of electrons and reduction of disorder in the system. Therefore, the magnetic frustration and orbital degeneracy are hard to remove, and no magnetic ordering is observed in SrNiRu₅O₁₁. However, the spin and orbital degrees of freedom of Ni are both evident from the sizable magnetization and Curie-Weiss behavior. Therefore, the Ni spins may order as temperature is lowered below 1 K, for example, via long-range dipole interactions. In this situation, we expect the competition between magnetic correlations among Ni²⁺ magnetic moments and geometrical frustration on the kagome nets to play an essential role in the precise nature of the ground state of SrNiRu₅O₁₁. In summary, the geometrical frustration and antiferromagnetic correlations within the Ru³⁺ ($S = 1/2$) triangular lattice result in very large ground state degeneracy with no long-range order at

any temperature, while the remaining interactions between the Ni²⁺ spins tend to order the system ferromagnetically. If these two processes are about equal in strength, then SrNiRu₅O₁₁ might lie near a spin-liquid phase transition.

IV. CONCLUSION

We observe a highly unusual combination of magnetic, transport, and thermal behaviors of single-crystal SrNiRu₅O₁₁. We speculate that an important role in the formation of this unusual ground state is played by competition between weakly interacting Ni²⁺ ($S = 1$) spins located between two-dimensional kagome layers, and geometrically frustrated Ru³⁺ ($S = 1/2$) spins located within the kagome nets. Although the compound does not show a clear signature of long-range magnetic order down to 0.35 K, residual magnetic Ni-Ni interactions induce a small nonlinear field dependence of magnetization at low temperatures. The Hall resistivity reveals no anomalous contribution down to 4.5 K, consistent with the absence of ferromagnetism. The power-law dependence of the electronic specific heat suggests a spin-liquid or TI state with massless Dirac fermions. The existence of exotic low-energy excitations at low temperatures might originate from the topological character of the tight-binding band structure of a kagome lattice [31].

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