Electronic in–plane symmetry breaking at field–tuned quantum criticality in CeRhIn₅

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Electronic nematic materials are characterized by a lowered symmetry of the electronic system compared to the underlying lattice, in analogy to the directional alignment without translational order in nematic liquid crystals. Such nematic phases appear in the copper- and iron-based high-temperature superconductors, and their role in establishing superconductivity remains an open question. Nematicity may take an active part, cooperating or competing with superconductivity, or may appear accidentally in such systems. Here we present experimental evidence for a phase of fluctuating nematic character in a heavy-fermion superconductor, CeRhIn₅ (ref. 5). We observe a magnetic-field-induced state in the vicinity of a field–tuned antiferromagnetic quantum critical point at \( H_c \approx 50 \) Tesla. This phase appears above an out-of-plane critical field \( H^* \approx 28 \) Tesla and is characterized by a substantial in-plane resistivity anisotropy in the presence of a small in-plane field component. The in-plane symmetry breaking has little apparent connection to the underlying lattice, as evidenced by the small magnitude of the magnetostriiction anomaly at \( H^* \). Furthermore, no anomalies appear in the magnetic torque, suggesting the absence of metamagnetism in this field range. The appearance of nematic behaviour in a prototypical heavy-fermion superconductor highlights the interrelation of nematicity and unconventional superconductivity, suggesting nematicity to be common among correlated materials.

Metals containing cerium (Ce) exhibit rich physics owing to their \( 4f \)² electrons, which can display a local character, acting as a magnetic moment, and can also hybridize with the conduction electrons leading to the formation of heavy quasiparticles, the so-called heavy fermions. Here we focus on CeRhIn₅, which at ambient pressure and zero magnetic field exhibits antiferromagnetic (AFM) order of nominally localized \( 4f \) electrons at Néel temperature \( T_N \approx 3.85 \) K. However, the notion of fully localized moments fails to capture its enhanced electronic specific heat \( C_e \approx 70 \) mJ mol⁻¹ K⁻²), suggesting that a small fraction of \( f \)-character persists at the Fermi surface even in the ordered state. Hydrostatic pressure tends to increase the \( -\gamma \)-orbital hybridization and thus generally leads to delocalization into heavy \( f \)-hybridized bands. This scenario seems to apply well to CeRhIn₅, where the critical pressure \( P_c \approx 23 \) kbar defines a quantum critical point (QCP) separating the low-pressure, local moment AFM phase from the high-pressure, delocalized paramagnetic phase at zero temperature. This can be directly observed by the pressure-induced change from a small to a large Fermi surface seen in quantum oscillation experiments, accompanied by a divergence of the quasiparticle mass around the QCP. A dome of unconventional superconductivity with a maximal transition temperature \( T_c \approx 2.2 \) K arises around the QCP with the related AFM quantum fluctuations being prime candidates for the pairing mechanism.

Similarly, at ambient pressure the AFM order can also be suppressed by strong magnetic fields at a critical field of \( H_c \approx 50 \) T. Unlike pressure, magnetic fields tend to localize the \( 4f \) electrons and thus the physical situation at the phase boundary is expected to be different. Yet the phenomenology in high magnetic fields and under pressure shares unexpected similarities. Recently a correlated phase was discovered in fields exceeding \( H^* \approx 28 \) T as a dome connected to the field-induced QCP at \( H_c \) (refs 8, 9), reminiscent of the dome of superconductivity under pressure.

Our main experimental observation is the appearance of strong in–plane resistivity anisotropy in this high-field phase beyond \( H^* \). Reliable, direction-dependent resistivity measurements under high magnetic fields are commonly difficult in highly conductive metals such as CeRhIn₅. To address this experimental obstacle, we have fabricated microstructured devices from a single crystal of CeRhIn₅ using focused ion beam (FIB) machining (details of the fabrication procedure are available elsewhere and are further described in the Methods section and Extended Data Fig. 1). Typical devices feature rectangular bars in the CeIn₃ plane (the \( a-b \) plane) of CeRhIn₅ (Fig. 1). The four-terminal transport bars were aligned along different crystal directions to probe different crystal symmetry channels. One microstructure design probes resistance in the plane along the [110], [100] directions, and a second along the symmetry–inequivalent [100], [010] directions. The zero-field resistances of all devices are in quantitative agreement with each other and with bulk measurements, excluding experimental artefacts due to strain or current path misorientation in the structures. The resistivity at room temperature is \( 31.3 \pm 1 \) μΩ cm at 2 K. The residual zero-field resistance ratio in excess of 100 indicates the high material quality of the crystal after the FIB process, as further supported by the observation of pronounced quantum oscillations.

Figure 2 encompasses the main experimental finding for both design types. The magnetoresistance of each device shows a resistive anomaly for fields around \( H^* \approx 28 \) T, signalling the entry into the high-field phase. At the same time, a substantial asymmetry between the orthogonal in–plane resistance bars emerges (red, blue traces in Fig. 2a). While one bar shows a strong increase in resistance, in the other the resistance decreases. A small in-plane field \( (H_{\parallel}) \) is key to controlling the directions of lower and higher resistance within the high-field phase. Here, the field is applied at a 20° angle with respect to the \( c \) axis, and tilted along one of the resistance bars.

The similarity of the anisotropy between the two inequivalent symmetry channels \( B_{1g} ([100],[010]) \) and \( B_{2g} ([110],[110]) \) is remarkable (Fig. 2). To simplify the comparison between geometries and samples, we define the anisotropy as the ratio of the resistivity along the in-plane field to the resistivity perpendicular to it, \( \rho_{\parallel}/\rho_{\perp} \). For fields \( 20° \) off the \( c \) axis, the in-plane anisotropy exceeds 5 at \( H^* \). It successively decreases at higher magnetic fields and becomes isotropic at the AFM phase boundary (Fig. 3). Although the shape of the transition is well reproduced among all five studied samples, the magnitude of the anisotropy appears to be sample-dependent, with peak values ranging between 4 and 10.

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When \( H_{\text{ip}} \) is rotated from one leg of the device towards the other, the resistivity anisotropy reverses and perfectly mirrors the previous measurement. At the same time, a small anisotropy is observed even at fields below \( H^* \). This small anisotropy is the result of a difference between the longitudinal (\( I \parallel H_{\text{ip}} \)) and transverse (\( I \perp H_{\text{ip}} \)) magnetoresistance in the presence of the in-plane field, which vanishes with \( H_{\text{ip}} \) as the field is turned closer to the \( c \) axis, as expected. The temperature dependence of the resistivity anisotropy is shown in Fig. 3b. The onset field is almost temperature-independent, similar to the pressure-induced transition into a non-magnetic superconducting state\(^{11} \). With increasing temperature, the magnitude of the anisotropy gradually shrinks until no anomaly is observed at \( H^* \) for temperatures above \( T^* \approx 2.2 \) K, which coincides with the maximal superconducting \( T_c \) under pressure. At higher fields the anisotropy gradually vanishes, and this upper field scale traces the field dependence of the Néel line, \( T_N(H) \).

The transport anisotropy appears to develop exclusively within the plane. Whereas the in-plane resistance is characterized by a large jump, the resistance along the \( c \) direction remains featureless (Fig. 2c). Despite the presence of strongly 3D electronic bands in CeRhIn\(_5\) (ref. 7), the in-plane resistivity anisotropy appears to occur in a highly 2D electronic system, which does not influence electronic transport along the \( c \) direction. A natural origin of such a resistance anomaly in a spin-spiral antiferromagnet\(^{12} \) could be a metamagnetic transition. Although metamagnetism in CeRhIn\(_5\) is well documented at low fields of 2 T (ref. 13), pulsed field magnetization up to the AFM phase boundary\(^{14} \)—as well as our own magnetic torque measurements under the field configuration where the anisotropy is most pronounced (\( \theta = 20^\circ \))—does not show any signature of a change in the susceptibility or magnetization at \( H^* \) (Fig. 2b; see Methods section for a detailed discussion of the metamagnetism in CeRhIn\(_5\)). Therefore metamagnetic transitions appear unlikely as the origin of the resistive anomaly.

The quantitative similarity between two symmetry-inequivalent channels \( B_{1g} \) and \( B_{2g} \) in the tetragonal crystal appears striking. This, as well as the increase of differential resistance with increasing current bias (Fig. 2e), is difficult to reconcile with a finite momentum order such as density waves as suggested by initial reports\(^{8,9} \) (see Methods). Alternative scenarios thus probably involve electronically broken in-plane symmetry with zero momentum transfer, that is, an electronic nematic state\(^1 \). In analogy to liquid-crystal phases in fluids that break rotational symmetry while preserving translational symmetries, the ‘electronic nematic’ describes a correlated state of matter that breaks a rotational symmetry of the underlying Hamiltonian while preserving translational symmetries. Only a few materials have been proposed to be electronic nematics, such as the 2D metal Sr\(_3\)Ru\(_2\)O\(_7\), (Al,Ga)As heterostructures and high-temperature superconductors. The observation of field-induced in-plane resistivity anisotropy in Sr\(_3\)Ru\(_2\)O\(_7\) has been a key experiment\(^{15} \) indicating its nematic character, and both its temperature and field dependence is strongly reminiscent of the anisotropy in CeRhIn\(_5\), depicted in Fig. 3. While fields along the \( c \) direction induce the transition in both compounds, a small in-plane field component \( H_{\text{ip}} \) selects the planar direction of high resistance. The in-plane fields required to align the nematic state are remarkably similar in these two compounds (CeRhIn\(_5\), 20°; Sr\(_3\)Ru\(_2\)O\(_7\), 13° off \( c \) axis\(^{15} \)), yet much larger in-plane fields are required in two-dimensional-electron-gas (2DEG) systems\(^{16} \). But whereas in Sr\(_3\)Ru\(_2\)O\(_7\), an increase in the resistivity is observed upon entering the nematic state for both orthogonal in-plane directions, in CeRhIn\(_5\) the resistivity increases in one direction and decreases in the other direction. Such behaviour is observed in the electronic nematic phase in the fractional quantum Hall state\(^{17} \). Intriguingly, the nematicity in all these previous examples appears to be strongly locked to the crystal, leading to an Ising-like character. In contrast, the high-field state of CeRhIn\(_5\) appears not to be strongly pinned to the lattice and hence would be a candidate for a new XY-nematic.

Figure 1 | CeRhIn\(_5\) microstructured devices for measurements of in-plane resistivity anisotropy. a, False-colour SEM image of a CeRhIn\(_5\) crystalline microstructured device. The crystal lamella (purple) was carefully aligned so that the four-point resistance bars are along the [110] and [100] directions, shown by blue and red arrows, respectively. All bars have identical dimensions, with a length of 20 \( \mu \)m, width of 4 \( \mu \)m and thickness of 1.8 \( \mu \)m. The green arrow \( H \) depicts the field configuration.

b, Similar microstructured device with resistance bars along [100] and [010]. c, Resistivity of the devices in \( b \) as a function of temperature. Both independent [100] legs of the device shown in \( b \) are given (left and right). The inset shows a zoom of the data about the AFM transition at 3.85 K.
Pure electronic nematicity is an idealized concept, as nematicity in real materials always appears intertwined with other order parameters. Symmetry dictates that once the rotational symmetry is broken by the nematic order, any other order parameter belonging to the same broken symmetry is bound to appear. For example, recent neutron diffraction experiments in Sr$_2$RuO$_4$ found evidence for a spin-density wave. Thus, in strongly interacting electron systems, electronic nematicity is inseparably entangled with lattice distortions or spin-textures, such as in the iron-based superconductors. This may also be the case in CeRhIn$_5$. A careful investigation of the high-field magnetization of CeRhIn$_5$ (figure 7 in ref. 14) suggests the presence of a minute change of magnetic susceptibility in fields around 30 T, which may be a signature of the spin texture relaxing into the new symmetry of the electronic system. Similarly to the magnetic structure, the lattice is expected to respond to the formation of electronic nematicity by breaking the four-fold symmetry of the underlying lattice. To probe the lattice response to the nematic transition directly, we performed fibre-optic dilatometry on CeRhIn$_5$ under pulsed magnetic fields tilted 11° away from the c direction. The data are consistent with a small lattice response just above the noise floor, with an upper limit of $\Delta L/L(H^*) < 2 \times 10^{-6}$ (see Methods section and Extended Data Fig. 4). Such a lattice distortion is at least two orders of magnitude smaller than the one observed at the field-induced transition in Sr$_2$RuO$_4$ (ref. 20), indicating that the lattice plays only a minor role in the anisotropy observed in CeRhIn$_5$. Thus, while the phase transition in CeRhIn$_5$ at $H^*$ may be reflected in subtle relaxations of the magnetism and the lattice, it appears to be dominated by the itinerant electronic system, which is a key characteristic of electronic nematics.

The absence of spontaneous macroscopic symmetry breaking in the presence of finite disorder is expected for XY-nematics in dimensions $D < 4$, and indeed no anisotropy at $H^*$ is observed for fields well aligned along the c direction (Fig. 2d). Two plausible scenarios could describe the presented data: (a) static, long range nematic order fragmented into domains; or (b) strong nematic fluctuations without static order. In light of the experimental limitations at low temperatures and high magnetic fields, these scenarios are difficult to distinguish conclusively. One experimental approach is to cycle the in-plane field component and search for hysteretic behaviour. If static domains are oriented by the application of in-plane fields, the oriented domain pattern may remain pinned when removing the in-plane field, leading to a history-dependent domain structure and a hysteresis loop. On the other hand, in the fluctuation scenario a hysteresis is naturally absent. This was probed by carefully aligning the field along the c direction, followed by increasing the field to 35 T (well above $H^*$) to induce a virgin domain pattern. Then the sample was rotated to 20°, where the anisotropy is maximal, and back to the c direction. No hysteresis is observed in this procedure (Fig. 2d). Nonlinear conductivity measurements can provide a similar test of domain reorientation. The appearance of nonlinear conductivity suggests that electric fields are also effective at aligning the nematic phase (Fig. 2e). The presence of domains can be probed by applying an electric field strong enough to orient them, followed by removing the d.c. bias and checking for remnant orientation. The differential resistance is fully reversible in current and no evidence for hysteresis has been observed.

The absence of hysteresis would be a natural consequence of a fluctuating nematic without static long-ranged order. However, CeRhIn$_5$ appears to be of XY-like nematic character, which unlike Ising- or Potts-like static orders is characterized by extended domain walls. XY-like domain walls are generally more weakly pinned owing to their smaller energy density, and thus it is conceivable that a reversible orientation of static nematic domains occurs. It thus remains a challenge for future studies to identify the microscopic nature of the electronic nematic. Certainly, the Ce 4f electrons and their interactions play an essential part in this phenomenon. One possible scenario would be a partial melting of the antiferromagnetically aligned in-plane spins into correlated 1D spin-chains driven by the quantum fluctuations of the

**Figure 2 | Broken tetragonal symmetry in the high-field state.**

**a.** Magnetoresistance for the two microstructured devices shown in Fig. 1 at 500 mK with fields tilted 20° away from the c axis. Both rows correspond to the design sketched on the left, and the colours of the resistance traces correspond to the bars of each device as indicated by the arrow colours. The top (bottom) row probes the B$_{xy}$ (B$_{xz}$) symmetry channel. The left (right) column has the in-plane component of the field aligned along the red (blue) current bar as indicated in the sketch at the top of the figure. **b.** Magnetic torque at 380 mK. The torque is featureless except for de Haas-van Alphen oscillations, suggesting the absence of metamagnetic transitions in high fields and in particular around $H^*$. The inset is a zoom of the data near $H^*$. The curves are labelled by the angle (in degrees) that the field was applied away from the c axis. **c.** In-plane and out-of-plane resistivity across the transition in a microstructure at 500 mK, for fields 20° away from the c direction. The high-field transition is indicated by the jump in the in-plane resistivity, yet the out of plane resistivity remains featureless. This measurement was performed on a device similar to those shown in Fig. 1 probing the (a, c)-anisotropy, which featured two independent c-axis resistance bars ($\rho_{ac}$, $\rho_{ac}$). **d.** Angular dependence of the in-plane resistivity along [110] (blue trace) and [100] (red trace) as the field is tilted towards the [110] direction. The anisotropy persists down to the lowest angles and gradually closes at 0°. The angle-down sweeps (turquoise, magenta) perfectly overlap with the up-sweep, indicating that no hysteresis can be detected. **e.** Nonlinear conductivity in the high-field state. Labels at right show levels of (strong) d.c. bias superimposed on the (small) a.c. (100 μA), thus probing the differential resistance in the presence of strong electric fields. Although the low-field state is perfectly ohmic, a pronounced nonlinear conductivity appears in the high-field phase.
The suppression of the AFM order under magnetic fields along the [110] direction (red trace), the anisotropy is exactly inverted compared to the case of a tilt along the [110] direction (blue trace). Figure 4 contrasts the AFM phase boundary under pressure and magnetic field, plotted against the dimensionless tuning parameter \( g/g_c = H/H_c \) or \( p/p_c \), respectively. The nematic phase bears remarkable similarities to the superconductivity observed under pressure: (1) both are bounded by a vertical line of first-order phase transitions; (2) they share the same temperature scale; and (3) both exist around the field- or pressure-driven destruction of the AFM order at similar values of \( g/g_c \).

In light of these similarities, a common microscopic mechanism would be an interesting way to unify the experimental observations, and we sketch our suggestion in Fig. 4b. Approaching the AFM phase boundary at low fields under pressure induces a region of d-wave superconductivity, whereas approaching it under large magnetic fields leads to an electronic nematic phase. Another intriguing aspect of this apparent nematicity in CeRhIn5 is the role of low electronic dimensionality. All previous candidate systems for electronic nematic materials such as Sr3Ru2O7, copper oxides, iron-based superconductors and 2DEGs heterostructures or La1.75Sr0.25NiO4 (ref. 28). Thus, CeRhIn5 is a particularly intriguing nematic candidate in which to experimentally investigate their relation, as both unconventional superconductivity and the nematic state can be accessed via the cleanest possible tuning parameters, namely, pressure and magnetic field. A direct comparison of the destruction of the AFM order by pressure and field indeed reveals unexpected parallels between the superconducting and the nematic state.
are generally characterized by an electronically layered structure. In contrast, CeRhIn₅ is a strongly 3D metal, as indicated by its low resistivity anisotropy of 1.8 (Fig. 2c) and quantum oscillation experiments which reveal Fermi surfaces of 3D character. Yet, despite this 3D structure, the nematic phase transition seems to occur in a strongly 2D electronic sub-system that dominates the in-plane transport but contributes negligibly to the out-of-plane conduction. This apparent 2D character again uncovers links between nematicity and superconductivity in CeRhIn₅, where detailed resistance measurements under pressure also indicate an unexpected 2D character of the superconducting order. Together, these findings highlight once more the importance of low-dimensionality for electronic nematic phenomena and their entanglement with unconventional superconductivity, even where it is not expected—such as in strongly 3D metals like CeRhIn₅.

Online Content Methods, along with any additional Extended Data display items and Source Data, are available in the online version of the paper; references unique to these sections appear only in the online paper.

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Author Contributions P.J.W.M. and F.R. designed the experiment. P.J.W.M., K.R.S., T.H. and M.D.B. fabricated the microstructured devices. P.J.W.M., T.H., M.K.C., B.J.R., R.D.M. and F.F.B. performed the pulsed field experiments and P.J.W.M., K.R.S. and L.B. the dc-field experiments. E.D.B. and F.R. grew the single crystals. M.J. performed the magnetostriiction measurements and their analysis. All authors contributed to the manuscript.

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Crystal growth and microstructuring. Single crystals of CeRhIn5 were prepared using indium flux. The samples were confirmed to have the tetragonal HoCoGa5 structure by X-ray diffraction measurements and were screened by resistivity and susceptibility measurements, which showed no detectable free indium. The high quality of the samples is reflected by the very small residual resistivity and the presence of quantum oscillations in transport and thermodynamics.

All devices used in this study were fabricated from one millimetre-sized single crystal of the AFM phase. The crystal was aligned by Laue diffraction, which agreed with the clearly visible tetragonal morphology of the crystal. The procedure for FIB microstructuring consists of four main steps: (a) lamella cutting; (b) undercutting and polishing; (c) ex situ extraction and transfer to the substrate, and (d) final structuring. The fabrication steps are related to FIB cutting of transmission electron microscope (TEM) lamella and have been described earlier (see, for example, ref. 9), and are sketched in Extended Data Fig. 1.

Lamella cutting. The first step consists of the fabrication of a typically 100 μm long, 40 μm wide and 3–4 μm thick oriented crystal slab, the lamella. Such a lamella is a typical starting material for a standard FIB microstructure fabrication. After orientation, the lamella is cut out using the ion beam in two steps. With a relatively large current of 9.3 nA, two trenches are cut into the material thereby producing the lamella from the crystal bulk.

Undercutting and polishing. The lamella is then undercut at an angle of 60°, to free it from the main crystal and facilitate the later extraction. In addition, a finer polishing step is applied to the sidewalls. This step leads to a cleaner surface for subsequent contacting and ensures a rectangular cross-section of the lamella. This step is performed at 2.5 nA and 30 kV. These currents and voltages are larger than those usually used in TEM lamella polishing due to the much larger size of the polished lamella. At the same time, we have not observed improvements in contacting or bulk transport properties for CeRhIn5 devices polished at lower currents or voltages.

Ex situ extraction and transfer to the substrate. The lamella is then extracted from the trench ex-situ under an optical microscope and transferred onto a Si/SiO2 substrate with predefined gold contacts. A small epoxy droplet is deposited on the substrate and the lamella is carefully placed on top of it. In suitably small droplets, the capillary forces ensure a flat, parallel and square mounting of the sample with respect to the substrate, as can be seen in Fig. 1. In a subsequent step, the device is sputter-coated with 200 nm of gold through a shadow mask for electrical contact. Sputter-depositing gold in an Ar-plasma chamber was found to provide superb contacts with sub-ohm contact resistances even for the 4 μm2-sized contact areas on lamellae of CeRhIn5.

Final structuring. The thus prepared device is then reintroduced into the FIB chamber for final structuring. At this stage, the rectangular lamella is patterned into the desired shapes to tailor the current path to the respective experiments, such as the [100],[101] or the [110],[110] in-plane anisotropy configurations shown in Fig. 1. During the final structuring, the thermal situation is distinctively different from the lamella coarse cutting described in (a) where the lamella is in good direct thermal contact with the large single crystal. The lamella mass and hence its total specific heat is small, and at the same time it is mounted by weakly heat-conducting epoxy glue such that the main path of heat removal in the vacuum environment of the FIB chamber is provided by the thin gold contact. Thus, it is important to minimize the beam power at this stage as much as feasible. The structures used in this study were fine-structured at a current of 800 pA.

Field induced quantum criticality. The suppression of the AFM order under pressure induces a dome of superconductivity around a QCP separating the AFM phase at low pressure from the paramagnetic phase characterized by delocalized f electrons at high pressure. The associated quantum critical fluctuations are prime candidates to mediate the unconventional superconductivity in CeRhIn5. Figure 4 shows the phase diagram of the AFM phase under the different tuning parameters of magnetic field and pressure, and highlights the similarities between them. In particular, it shows the appearance of a correlated phase (superconductivity under pressure, the nematic state under field) across the vertical line of first-order phase transitions at similar relative values of the tuning parameters and up to a similar temperature scale of 2.2 K. Additional similarity arises when examining the temperature dependence of the resistivity.

As the AFM phase boundary delineates a line of broken symmetry, it must be accompanied by a phase transition at zero temperature at \( T_c \approx 50 \text{ K} \). The nature of the critical fluctuations, however, may be expected to be different. Under pressure at \( p_c \), the f electrons are delocalized and the entanglement fluctuations of the unbroken moments lead to the formation of correlated states, such as unconventional superconductivity. In a magnetic field, the situation may be different as fields generally increase the localization of the f electrons and the resulting high-field state may be a field-polarized paramagnetic state with weak fluctuations. Thus, a question about the quantum critical nature of the transition at \( H_c \) arises.

Our resistivity measurements are compatible with a QCP at \( H_c \). The longitudinal magnetoresistance at multiple temperatures for \( H \) is shown in Extended Data Fig. 2. From these data, the temperature dependence for various fields can be extracted. For fields below the critical field \( H_c \approx 49.7 \text{ T} \), the resistivity shows a drop at the AFM phase boundary. At higher fields the low-temperature resistivity crosses over into a \( T^2 \) Fermi-liquid-like behaviour. At \( H_c \), the resistivity is sublinear with a power-law behaviour, \( \rho_{||}(H) \approx 0.75 \mu \Omega \cdot \text{cm} \times T^{-0.35} \). This power law is nearly identical to the zero-field temperature-dependent resistivity at a field \( H \). This can be typically seen in organic 1D conductors; and (b) the loss of scattering and opening of a partial gap on the Fermi surface leads to a resistance increase, which is its apparent XY-character as discussed in the main text. This is unexpected for finite-\( q \) orders, in which the directionality of the order arises along different \( q \) vectors connected to the electronic instability. In the case of density waves, the direction is given by the Fermi surface morphology (often simplified as Fermi surface nesting) and the electronic susceptibility.

The observations summarized in Fig. 2 suggest that the in-plane field component selects the direction of reduced resistivity, and that this resistivity directionality can be equally directed along the [100] and [110] directions.

Here we show that fine angle scans around \( H \) confirm and support this equivalence. Extended Data Fig. 3 contrasts the two situations of in-plane fields along [100] and along [110]. The quantitative agreement between these two directions is striking (dashed versus solid line). The resistivity was calculated only from the geometric dimensions measured by electron microscopy. The similar resistivity and angle dependence of the anisotropy indicates that the lattice plays only a minor role in determining the direction of the scattering, which appear to be set entirely by the direction of the in-plane field.

In discussing its nematic character, it is important to note that the two experiments determining the anisotropy of [100],[101] and [110],[110] probe symmetry-invariant directions. In the tetragonal symmetry group, these probe quantities belonging to different irreducible representations, \( B_{2g} \) and \( B_{1g} \), respectively. Therefore, there is no symmetry-related reason to expect such a quantitative similarity between these experiments. In particular, this is inconsistent with the assumption of stripe-type orders as observed in the copper oxides. There, a clear alignment of 1D conduction channels through the formation of stripes was suggested to explain the appearance of an in-plane resistance anisotropy. These striped lines are found to run along the nearest neighbours and thus induce a [100],[010] anisotropy. In the experimental \( B_{2g} \) configuration, [110],[101] such order would not lead to apparent anisotropy.

Nematics and density waves. Here we discuss in more detail the nature of nematic order in contrast to spin- and charge-density-waves. The key characteristic of an electronic nematic is preserving the translational symmetry (\( q = 0 \)). Density waves at finite-\( q \) vectors would also be natural candidates for a high-field order that breaks the rotational symmetry and would lead to transport anisotropy. We find our recent observations to be inconsistent with the expectations for density waves. However, the resistivity and transport measurements point to a picture different from conventional density wave orders, microscopic techniques that directly probe the momentum dependence of the high-field order would be desirable. Absence of resistive anomaly for \( H \) \( |c| \). A transition into a density wave state is generally signalled by a sudden change of resistivity qualitatively similar to our observation in Fig. 2a. Two mechanisms compete in determining the resistive signature of density wave transitions: (a) the decrease of charge carriers due to the opening of a partial gap on the Fermi surface leads to a resistance increase, which can be typically seen in organic 1D conductors; and (b) the loss of scattering channels as the Fermi surface is reconstructed into smaller pockets reduces the resistivity. The image of 2D and 1D materials—for example, spherical to ellipsoidal through a Pomeranchuk instability—is its apparent XY-character as discussed in the main text. This transition to lower order decreases the Fermi velocity in one direction and increases it in the orthogonal direction, with a velocity anisotropy linear in the nematic order parameter. This is distinct from Ising-like nematic orders, for example in 2DEG systems where the nematic director is strongly pinned to the [110] direction.
Absence of angular hysteresis of the anisotropy. The disappearance of anisotropy for \( H_c \) can only be explained in the context of density waves or charge orders by assuming a microscopic arrangement of randomly oriented domains which mask the microscopically broken \( C_3 \) symmetry on the macroscopic scale. Given the small overall size of the investigated microdevices, such domains would need to be in the nanometre range, which is not uncommon for such orders in correlated metals. In this picture, the in-plane field component would act to align the domains. Thus, if a state is prepared at \( H_c \) and completely fractured into nanodomains, it could not be possible to show no anisotropy (see previous section). Then, as the in-plane field is increased, the domains react by realigning themselves with the in-plane field, energetically favourable patches grow in size and unfavourable ones shrink. As charge- and spin-density-waves are finite-\( q \) orders, they can couple effectively to defect potentials in the crystal of length scale \( \sim 1/q \). Such defects are effective pinning centres that fix the phase of the density wave. In the experimentally unobservable presence of pinning, one would thus expect a hysteresis due to remnant alignment when the in-plane field component is brought back to zero. Figure 2d shows this experiment: first the field was increased for \( H = 35 \) T, where no anisotropy was observed. Then it was rotated into static field to 20°, which marks the maximum of the anisotropy and thus could be taken as an indication for the maximal alignment of density wave domains. Then the field was rotated back to \( H_c \), and the data overlap within experimental accuracy without any sign of a hysteresis loop. The smaller features are not experimental noise but reproducible finite structure most probably related to the Fermi surface morphology, similar to angle-dependent magneto-resistance oscillations (AMROs). The absence of hysteresis would be well compatible, however, with either nematic fluctuations or static \( XY \)-like nematic domains. \( XY \)-like order parameters are known to have extended domain walls minimizing the line energy associated with such walls. Therefore, the polarization process gradually aligning the nematic directors along the in-plane field leads naturally to less hysteretic and history-dependent behaviour.

Nonlinear conductivity. The transition into the high-field phase is not detectable by in-plane resistance measurements when the field is well aligned with the \( c \) direction (Fig. 2e). To probe this state in more detail, we have performed nonlinear conductivity measurements obtained by superposition of a d.c. bias on top of a small a.c. excitation \( I_{ac} = 100 \mu A \). Whereas in the absence of d.c. bias, no signature of the transition appears for purely out-of-plane fields, the transition is uncovered by a strong nonlinear increase of the resistance upon increasing bias. This is contrary to the expectations for density wave materials. Here, the gapped quasiparticles that participate in the density wave start to contribute to the charge transport above a threshold electric field due to the onset of sliding motion of the condensate. This well-known depinning of density waves leads to characteristic nonlinear conductivity that increases with increasing electric field as an additional conduction channel appears under high bias. The measurement presented in Fig. 2e shows the opposite behaviour. As the electric field is increased, the conductivity decreases in this configuration. While it cannot be explained by density wave motion, it would be consistent with in-plane electric fields aligning nematic domains akin to the role of the in-plane field component.

In a previous study, we have observed a negative differential conductance in the presence of a perpendicular in-plane field at 20° away from the \( c \) direction (the maximal resistance state). This suggests a competition between the electric and magnetic in-plane fields in determining the alignment of the high-field order. Further detailed studies of the angle dependence under high electric field bias are required to uncover this combined effect. However, it is worth noticing that the observation presented here of negative differential conductivity in any magnetic and electric field combination is challenging to explain in the framework of density waves.

In agreement with the absence of hysteresis as a function of angle, we similarly also do not observe a hysteresis in current. The d.c. bias can be increased and returned to zero without any detectable change at zero bias current. This absence of hysteresis and history-dependence is at odds with similar arguments of domain reorientation that explain the various experimental reports on hysteretic and history-dependent transport phenomena in density waves, often named ‘memory effect’.

Role of the lattice. To investigate the role of the lattice at the \( H_\ast \) transition, we measured the magnetostriction along the \( c \) direction in pulsed fields using an optical fibre Bragg grating (FBG) technique. The experiment was carried out using a 2-mm-long FBG sensor inscribed on a telecom-type 125-μm optical fibre, attached to the sample with acrylic-based glue. Broadband light in the infrared spectrum from a superluminescent light emitting diode (SLED) source injected into the fibre is reflected by the FBG at \( \sim 1.550 \) nm, then dispersed with a monochromator and acquired with a 50 kHz InGaAs linescan camera. This technique has previously been used to uncover magnetostriction signatures of similar and larger magnitude in other systems, associated with microscopic mechanisms in frustrated magnetic systems, multiferroics and metals that exhibit field-induced Lifshitz transitions.

Extended Data Fig. 4 shows the coefficient of thermal expansion \( \alpha(T) \) measured on a \( c \)-axis needle sample of CeRhIn\(_5\) of dimensions \( 0.5 \) mm \( \times 0.5 \) mm \( \times 2.8 \) mm during cooldown. A clear anomaly signalling the AFM transition is observed at \( T_\ast \), highlighting the sensitivity of the optical measurement technique.

The magnetic field pulses (namely to 32 T, 36 T, 40 T, 45 T and 55 T), thus minimizing the impact of artefacts originating in mechanical vibrations of the magnet. At the same time the observed anomaly is small, and further experiments in static magnetic fields are clearly required to quantify the lattice response at the potential nematic transition. Given the present data, however, we can exclude a strong field induced structural transition at \( H_\ast \) (upper limit \( \Delta L/L < 2 \times 10^{-5} \)). This is clear evidence that lattice response at \( H_\ast \) is at least two orders of magnitude smaller than the one observed in \( Sr_3Ru_2O_7 \) (ref. 20).

The reduced lateral dimensions of the probes used in our pulsed field experiment preclude bending of the fibre to obtain in-plane magnetostriiction for magnetic fields applied close to the \( c \) axis. All crystallographic axes are, nevertheless, correlated by thermodynamics and such correlation is quantified by Poisson’s ratio. A material such as CeRhIn\(_5\), displaying field-induced broken symmetry in the \( a-b \) plane, should consequently display a distortion along the \( c \) axis. Interestingly, this distortion should be evident even in the presence of in-plane domains, making the \( c \)-axis lattice parameter an ideal witness.

Metamagnetism in CeRhIn\(_5\). One prime candidate for a resistive anomaly under high magnetic fields in antiferromagnets characterized by complicated spin-patterns are metamagnetic spin-reorientations. These are quite common in Ce-bearing intermetallics, and it is thus important to understand the magnetic structure and its metamagnetic transitions.

In zero field, the AFM state of CeRhIn\(_5\) consists of anti-parallel planar spins on the Ce site forming an incommensurate spin-spiral along \( c \) with an ordering vector \( q = (1/2,1/2,0) \) (ref. 27). Indeed this structure exhibits metamagnetic transitions at low fields, as commonly found in incommensurate spin-spirals. At relatively low in-plane fields of \( 2 \) T, the spiral alignment of the spins along the \( c \) axis becomes energetically unfavourable and the moments align perpendicular to the in-plane component in a metamagnetic spin reorientation, locking into a commensurate order of \( q = (1/2,1/2,1/4) \). This low-field transition is observed as a sharp anomaly in the magnetization, magnetic susceptibility and specific heat.

Consistently, this metamagnetic spin reorientation is evident in our data as a step in the resistivity at a field of \( 5 \) T, which at \( \theta = 20° \) corresponds to \( \sim 2 \) T of in-plane field as expected from its \( \sin^2(\theta) \) angular dependence (Fig. 2f). This self-consistently shows that metamagnetic transitions can be detected by our methods, and that they occur in both resistive and magnetic measurements, unlike the transition at \( H_\ast \) which does not involve any apparent change in magnetism.

In fields beyond the metamagnetic transition, a gradual canting of the spins towards the eventual field-polarized paramagnetic state at \( 30 \) T has been proposed to satisfactorily explain the high-field magnetic behaviour up to the field-induced destruction of the AFM order. The absence of a step or change in slope of the longitudinal magnetization at \( H_\ast \) implies that the net magnetic moment of the spin structure parallel to a field applied along the \( c \) direction remains unaffected by the transition. At the same time, the magnetic torque, \( \tau = M \times H \), directly measures changes of the net magnetization of the spin structure perpendicular to the applied magnetic field. No anomaly in the torque appears, either at small angles close to \( H || c \) (the torque vanishes for fields exactly applied along crystal symmetry directions and thus cannot be measured exactly at \( H || c \)) or at the field angle of maximal resistivity anisotropy at \( \theta = 20° \). This absence of magnetic signatures puts clear constraints on the field dependence of the net magnetization, and effectively rules out any change in the fields used.

In principle, however, microscopic spin reorientations that change neither the net longitudinal nor transverse magnetization are possible between two complex AFM structures that accidentally produce the same net moment. This is quite exotic, as typically the Zeeman energy is the driving force of metamagnetic transitions, and without a change in the field-induced susceptibility at \( H_\ast \) no net Zeeman energy gain occurs.

It is also important to note that while our investigation focuses on the field configuration of maximal anisotropy (\( \theta = 20° \)), the anisotropy appears gradually as the field is rotated away from \( H || c \) (Fig. 2). This is contrary to the expec-
tions of a metamagnetic transition driven by the in-plane field, which would require a finite in-plane critical field. This agrees well with the experimental evidence for the occurrence of the transition at $H \parallel c$, as indicated by the stark change in the nonlinear transport (Fig. 2) as well as the experiments reported in ref. 9. At the same time, in the absence of an in-plane field, $H_{ip}$ neither an in-plane resistivity anisotropy between the orthogonal transport bars nor an isotropic feature shared by both in-plane directions appears. This would imply that at such a hypothesized metamagnetic transition at $H^*$, the spin scattering does not only remain isotropic but also quantitatively unchanged, and only the in-plane field component somehow influences the high-field spin state to show the observed tenfold scattering anisotropy.

Although it may be possible to construct a highly exotic metamagnetic reorientation around these strong constraints, it would necessarily involve a large degree of artificial fine-tuning. On the other hand, these properties naturally follow in a nematic picture as proposed in the main text.

**Data availability.** Source Data for Figs 1–4 are available in the online version of the paper. Other data generated or analysed during this study are available from the corresponding author on reasonable request.

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Extended Data Figure 1 | Sketch of the focused ion beam fabrication process. a. Precutting of the lamella from the large single crystal. Red arrow in all panels shows the position of the focused ion beam (FIB). b. Undercutting the lamella to disconnect it from the crystal for the \textit{ex situ} transfer. c. Transfer of the lamella onto a silicon substrate and electric contacting with gold (Au). d. Final fine structuring into the shapes shown in Fig. 1.
Extended Data Figure 2 | Quantum critical transport. a, In-plane resistivity as a function of field $H \parallel [010]$ for different temperatures (labels on curves, in K). The transition into the AFM state is clearly visible as a kink in the curve moving to higher fields at lower temperatures. At the lowest temperatures, the signature turns into a step-like transition.

b, Temperature dependence of in-plane resistivity at fixed field (see key: $H_c$, critical field) extracted from the data in a. Whereas the resistivity above and below $H_c$ cannot be described by a simple power law, at $H_c$ a slightly sublinear power law $\rho \approx T^{0.91}$ describes the measured resistivity over a temperature range up to 10 K.
Extended Data Figure 3  |  Angle dependence of the anisotropy. Main panel, measurements of the field angle dependence of in-plane resistivity in a static field (35 T applied along the c direction to ensure the sample is in the high-field phase) and at fixed temperature (300 mK). The samples were rotated in the field towards 20° at which the maximum anisotropy occurs. Two different samples were used to probe this anisotropy. The triangular one shown in Fig. 1 probes the anisotropy developing along the [110], [11̅0] channel (dashed line; also shown in Fig. 2), while a second similar sample featuring resistance bars along [100], [010] probed the other orientation (solid line). The emergent anisotropy of the two devices with similar dimensions is remarkably similar. The field configuration during the rotation is sketched on the right. The top sketch indicates the rotation of the magnetic field towards one in-plane direction (red arrow) while the orthogonal in-plane direction defines the axis of rotation (blue). The same colour-code is used consistently in the main panel. The middle (bottom) sketch indicates the configuration for the concrete case of the [110], [11̅0] ([100], [010]) device. Owing to the off-axis field during the rotation, an in-plane field $B_{ip}$ emerges and its direction in the individual devices is indicated by the green arrow.
Extended Data Figure 4 | Lattice response to temperature and magnetic field. a, Measured zero field coefficient of thermal expansion $\alpha(T)$ of CeRhIn$_5$ in $^3$He gas. The lambda-type anomaly at the second order AFM phase transition is evident ($T_N$, Néel temperature). b, Magnetostriction $\Delta L/L(H)$ measured along the c direction in pulsed magnetic fields at 510 mK. The pulsed field measurement was performed on the same needle-like crystal of CeRhIn$_5$ that is shown in the zero-field measurement in a. The high-field measurement was taken using the same experimental set-up as the temperature-dependent measurement. The crystal was immersed in liquid $^3$He with the magnetic field applied 11° off the c axis. These raw data show a typical smooth and roughly quadratic behaviour in the magnetic field. Both pulse upsweep (red) and downsweep (blue) overlap well and are barely distinguishable in the raw data. c, Same magnetostriction $\Delta L/L$ as in b but after subtraction of a smooth background. In the downsweep, a small but sharp feature appears at $H \approx 30$ T that may indicate c-axis shrinkage as the sample transitions into the nematic phase.