Unusual high-field metal in a Kondo insulator

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Strong electronic interactions in condensed-matter systems often lead to unusual quantum phases. One such phase occurs in the Kondo insulator YbB12, the insulating state of which exhibits phenomena that are characteristic of metals, such as magnetic quantum oscillations1, a gapless fermionic contribution to heat capacity2,3 and itinerant-fermion thermal transport4. To understand these phenomena, it is informative to study their evolution as the energy gap of the Kondo insulator state is closed by a large magnetic field. Here we show that clear quantum oscillations are observed in the resulting high-field metallic state in YbB12; this is despite it possessing relatively high resistivity, large effective masses and huge Kadowaki–Woods ratio, a combination that normally precludes quantum oscillations. Both quantum oscillation frequency and cyclotron mass display a strong field dependence. By tracking the Fermi surface area, we conclude that the same quasiparticle band gives rise to quantum oscillations in both insulating and metallic states. These data are understood most simply by using a two-fluid picture in which neutral quasiparticles—contributing little or nothing to charge transport—coexist with charged fermions. Our observations of the complex field-dependent behaviour of the fermion ensemble inhabiting YbB12 provide strong constraints for existing theoretical models.

In Kondo insulators (KIs), an energy gap is opened up by strong coupling between a lattice of localized moments and the extended electronic states. The resulting Kondo gap \( E_g \) is usually narrow (typically \( E_g \approx 5–20 \text{ meV} \)), yet the role it plays in charge transport is more complicated than that of the bandgap in conventional semiconductors. A low-temperature \( (T) \) saturation of resistivity has long been known in two prototypical KIs, namely, samarium hexaboride (SmB6) and ytterbium dodecaboride (YbB12) \( (\text{refs. 4–5}) \); both are mixed-valence compounds with strong \( f-d \) hybridization that defines the band structure close to the Fermi energy. While the saturation might suggest additional metallic conduction channels, the high resistivity value within the weakly \( T \)-dependent ‘plateau’ implies that these are highly unconventional\(^{14} \); one interpretation is the presence of topologically protected surface states\(^{6} \).

Recently, magnetic quantum oscillations, suggestive of a Fermi surface (FS) and thus totally unexpected in an insulator, have been detected in both SmB6 and YbB12 \( (\text{refs. 12–16}) \). Whilst some have attributed the oscillations in SmB6 to residual flux\(^{13} \), the flux-free growth process of YbB12 \( (\text{Methods}) \) excludes such a contribution. The oscillations in YbB12 are observed in both resistivity \( \rho \) (the Shubnikov–de Haas (SdH) effect) and magnetization \( M \) (the de Haas–van Alphen (dHvA) effect) at applied magnetic fields \( H \) where the gap is still finite. The \( T \) dependence of the oscillation amplitude follows the Lifshitz–Kosevich (LK) formula, which is based on the Fermi liquid theory\(^{7} \). Moreover, a contribution from gapless quasiparticle excitations to the heat capacity \( C \) has been detected in both KIs\(^{8,9} \). Though the thermal transport observations remain controversial in SmB6\(^{10,11} \), YbB12 shows a \( T \)-linear zero-field thermal conductivity \( \kappa \), a characteristic of itinerant fermions\(^{1} \). Agreement among the FS parameters derived from quantum oscillations, heat capacity and thermal conductivity in YbB12 suggests that the same quasiparticle band is responsible\(^{9} \).

Despite this apparent consistency, the mystery remains: how can itinerant fermions exist in a gapped insulator and transport heat but not charge? In response, many theoretical models entered the fray, including magnetoeectrons\(^{10} \), Majorana fermions\(^{13,16} \), emergent fractionalized quasiparticles\(^{17,18} \) and non-Hermitian states\(^{19} \). As these scenarios frequently envisage some form of exotic in-gap states, it is potentially invaluable to observe how the properties of KIs evolve as the energy gap \( E_g \) closes.

The cubic rare-earth compound YbB12 is an excellent platform to carry out such studies. In YbB12, the Kondo gap \( (E_g \approx 15 \text{ meV} \text{ (ref. 20)}) \) is closed by large \( H \), leading to an insulator-to-metal (I–M) transition at fields ranging from \( \mu_0H_{\text{K}} \approx 45–47 \text{ T} \) \( (\mu_0 \text{ is the vacuum permeability; field vector, } \text{H} \parallel \text{[100]} \) to 55–59 T \( (\text{H} \parallel \text{[110]} \text{ (ref. 1226)}) \). As revealed by a large Sommerfeld coefficient \( \gamma \) of the heat capacity\(^{2} \), Kondo correlation does not break down at the I–M transition, remaining strong to 60 T and beyond in the high-field metallic state; hence, this can be termed a Kondo metal (KM)\(^{21} \). In this study, we apply both transport and thermodynamic measurements, including \( \rho \), penetration depth, \( M \) and magnetostriiction, to YbB12. By resolving quantum oscillations and tracking their \( T \) and \( H \) dependence in the KM state, we trace the fate of the possible neutral quasiparticles at fields above the gap closure and expose their interactions with more conventional charged fermions.

In our YbB12 samples, \( M \) and magnetostriiction data show that the I–M transition occurs at \( \mu_0H_{\text{K}} \approx 46.3 \text{ T} \left( \text{H} \parallel \text{[100]} \right) \); the tiny valence increase in Yb ions at the transition suggested by the magnetostriiction reinforces the KM nature of the high-field metallic state \( \text{Methods} \) and \( \text{Extended Data Fig. 1} \). To probe the electronic structure of the KM state, a proximity detector oscillator \( \text{(PDO)} \) was used \( \text{Methods} \) for contactless SdH effect studies. The PDO technique is sensitive to the sample skin depth, providing a direct probe of changes in conductivity of the metallic KM state. The setup illustrated in the inset of Fig. 1a was rotated on a cryogenic goniometer to achieve \( \text{H} \)-orientation-dependent measurements. Fig. 1a
summarizes the $H$ dependence of the PDO frequency $f$ as $H$ rotates from $[100]$ to $[110]$.

The low conductivity of the KI state suggests that the megahertz oscillatory field from the PDO coil completely penetrates the sample\cite{supplementary}. Therefore, the response of the PDO is dominated by the sample skin depth only when the sample enters the KM state and the conductivity increases substantially (Methods). The transition between these regimes is marked by a dip in the $f$ versus $H$ curves close to $H_{\perp KM}$ above which magnetic quantum oscillations emerge. Figure 1b displays $\Delta f$, the oscillatory component of $f$, at various angles $\theta$ as a function of $1/H$. The distinct oscillation pattern observed for $H \parallel [100]$ ($\theta = 0$) is preserved up to $\theta \approx 21^\circ$ (Fig. 1b), being strongly modified at higher angles (Extended Data Fig. 2a).

The oscillations in $\Delta f$ represent a single series that is intrinsically aperiodic in $1/H$; attempts to fit them using a superposition of conventional oscillation frequencies fail to reproduce the raw data (Supplementary Information). To demonstrate the point further, Fig. 1c compares Landau-level indexing plots for the low-field KI and high-field KM states, both with $H \parallel [100]$. The oscillatory component of resistivity, $\Delta \rho$, in the KI state is shown in the inset of
Fig. 1c and conventionally indexed using integers for minima and half integers for maxima in \( \rho \). For the KM state, peaks in \( f \) correspond to peaks in conductivity\(^2\) and therefore are indexed using integers\(^2\). A further subdivision of the oscillations—reminiscent of a second harmonic—is likely due to Zeeman splitting of the quasiparticle levels\(^3\); these features are marked with ‘+’ and ‘-’, assuming the signs expected for conventional Zeeman shifts. In the KI state, the plot of Landau-level index \( N \) versus \( 1/\hbar \) is a straight line, as expected for a field-independent quantum oscillation frequency in a non-magnetic system. By contrast, in the KM state, the \( 1/\hbar \) positions of the oscillations have a nonlinear relationship with \( N \), which will be described below (here we note that the magnetic induction \( B \approx \mu_0 H \), since \( \mu_0 H \) is large and YbB\(_1\)_\(_3\) has weak magnetization; see Methods).

Nevertheless, despite their unusual periodicity, the \( T \) dependences of individual oscillation amplitudes in the KM state (Fig. 1d) closely follow the LK formula\(^2\)

\[
\Delta f(T) \propto \frac{2e^2 k_B T/\hbar}{\sinh(2\pi k_B T/\hbar^2)}.
\]

This suggests that they are almost certainly due to fermions. In this case, \( k_B \) is the Boltzmann constant and \( \hbar^2 \) is the cyclotron energy. However, the derived value of \( \hbar^2 \) varies nonlinearly with \( H \), indicating that the cyclotron mass \( m^* = eB/\hbar^2 \) is a function of \( H \) (Fig. 2c).

The relationship between oscillation indices and magnetic field is empirically described by

\[
N + \lambda = \frac{F_0}{\mu_0 (H_N - H^*)},
\]

where an offset field of \( \mu_0 H^* = 41.6 \text{T} \) has been subtracted from \( \mu_0 H_N \). Here \( \lambda \) is a phase factor, \( N \) is the index, \( H_N \) is the field at which the corresponding feature (for example, peak) occurs and \( F_0 \) is the slope. Equation (2) is symptomatic of an FS pocket that progressively depopulates as \( H \) increases, for reasons discussed below. In such cases, the Onsager relationship \( F(B) = \frac{h}{\pi} A(B) \) between the FS extremal cross-sectional area \( A \) and the frequency \( F \) of the corresponding quantum oscillations still applies even when \( A \) changes with \( H \) (Methods). Since \( B \approx \mu_0 H \), we write \( B^* = \mu_0 H^* \) and represent the field dependence of equation (2) using a \( B \)-dependent frequency:

\[
F_{KM}(B) = \frac{F_0}{B} \ln \frac{B}{B-B^*};
\]

this is associated with a \( B \)-dependent extremal area \( A(B) = \frac{2e^2}{\pi h} B \), where \( A_0 = \frac{2e^2}{\pi h} F_0 \). We mention that these results are consistent with the ‘back-projection’ approach\(^2\) for the measured SdH frequency (Supplementary Information).

Analysed in these terms, our data indicate that the quantum oscillations seen in the KM state are due to an FS pocket that is as far as very closely related to—the FS pocket in the KI state, which contributes quantum oscillations and a \( T \)-linear term in both \( C \) and \( \kappa_\theta \) (refs. \(^2\), \(^4\)). This is strongly suggested by the field dependence of \( F_{KM} \) (equation (3)), as shown in Fig. 2b. Even a cursory inspection reveals that equation (3), describing oscillations in the KM state, gives a frequency close to that of the KI state oscillations when extrapolated back to \( H_{KI} \). This is true for all angles \( \theta \) at which the oscillations were measured; using appropriate \( F_0 \) values (Supplementary Table 1) and substituting fields \( B = \mu_0 H_{KI}(\theta) \) on the phase boundary (minima in the PDO data; Fig. 1a) into equation (3) gives the frequencies shown as magenta diamonds in the inset of Fig. 2b. The \( \theta \) dependence of \( F_{KM}(\mu_0 H_{KI}) \) thus deduced tracks the behaviour of the dHvA frequencies measured in the KI state, albeit with an offset of \( \pm 100 \text{T} \) (Fig. 2b). This offset may be due to a discontinuous change in \( F \) at the phase boundary; however, it could also result from the potential uncertainty in determining the \( H \) position of any phase transition associated with a valence change\(^2\). Substituting \( \mu_0 H_{KI} \) values decreased by \( \sim 0.8 \text{T} \) into equation (3) yields an exact match of \( F_{KM} \) with the quantum oscillation frequencies observed in the KI state (Fig. 2b, inset).

Further evidence that the same FS pocket contributes to quantum oscillations in both KI and KM states comes from the cyclotron masses. As shown in Fig. 2c, a cyclotron mass \( m^* \approx 7m_e \) (\( m_e \) is the free electron mass) is deduced from the \( T \) dependence of the quantum oscillations at fields just above and just below the \( 1\)-\( M \) transition (\( m^* \) in the KI state is obtained from the dHvA oscillations, which we suggest are more fundamental; see Methods). As mentioned above, \( m^* \) in the KM state is enhanced by increasing the field. This effect is also captured by the unusual field-dependent Zeeman splitting of the SdH peaks (Extended Data Fig. 2, inset). In Fig. 2d, we plot the spin-splitting factor \( S \) (ref. \(^2\)) calculated from the SdH data (Methods). The field enhancement of \( S \) generally tracks that of \( m^* \), consistent with elementary expectations of Landau quantization (Methods). The scaling yields a \( g \) factor of \( \sim 0.084 \), much smaller than that for typical weakly interacting electron systems (\( g \approx 2 \)) (ref. \(^2\)).

The consistency of the frequency and mass across the \( 1\)-\( M \) phase boundary indicates that the novel quasiparticles detected in the KI state of YbB\(_1\)_\(_3\), which are probably charge neutral\(^1\), also cause the SdH effect in the KM state (see Methods for further discussion). The unusual nature of the KM state is further revealed by magnetotransport experiments. To reduce Joule heating as the KM state is traversed, we used a pulsed-current technique (Methods). Current is only applied to the sample when \( H > H_{KI} \), as shown in the inset of Fig. 3a. Below \( 10 \text{K} \), very weak longitudinal magnetoresistance (MR) is observed above \( 50 \text{T} \) (Extended Data Fig. 3) and is preserved up to \( \sim 68 \text{T} \) (Extended Data Fig. 2), permitting an analysis of the \( T \) dependence of resistivity \( \rho \) in the KM state. Figure 3a,b shows \( \rho \) at \( 55 \text{K} \) as a function of \( T \) and \( T^2 \), respectively. The \( \rho - T \) curve shows a maximum at \( T_* = 14 \text{K} \). With decreasing \( T \), a linear \( T \) dependence develops below \( 9 \text{K} \) and extends down to \( T^2 \approx 4 \text{K} \); subsequently, a Fermi-liquid-like \( T^2 \) behaviour is established below \( T_\text{FL} = 2.2 \text{K} \). The overall behaviour of \( \rho(T) \) mimics that of a typical Kondo lattice, where Kondo coherence develops at \( T^* \) and a heavy Fermi liquid state forms below \( T_\text{FL} \) (ref. \(^{30}\)).

The residual resistivity \( \rho(T \rightarrow 0) \approx 0.34 \mu\Omega \cdot \text{cm} \) indicates that the KM state may be classified as a ‘bad metal’.

Owing to the relatively high resistivity in the KM state, the Kadowaki–Woods (KW) ratio \( A_\theta /\gamma^2 \) (\( A_\theta \) is the \( T \)-coefficient of resistivity) is surprisingly large. Using the value of \( A_\theta \) obtained from the fit in Fig. 3b and the Sommerfeld coefficient given by the pulsed-field heat capacity measurements (\( \gamma \approx 63 \text{mJ mol}^{-1} \text{K}^{-2} \) at \( 55 \text{T} \) (ref. \(^1\); see Methods), the estimated KW ratio is \( 1.46 \times 10^3 \mu\Omega \cdot \text{cm} \text{K}^{-2} \text{mol}^{-1} \)), three and four orders of magnitude larger than typical values for heavy-fermion compounds and transition metals, respectively (Fig. 3c). Such an abnormal KW ratio cannot be addressed by the degree of degeneracy of the quasiparticles, which tends to suppress the KW ratio in many Yb-based systems\(^{30,35}\).

It should be clear that the fermions responsible for the high-field quantum oscillations cannot, on their own, account for the transport and thermal properties of the KM state. First, even if they were capable of carrying charge, their field-dependent SdH frequency and mass would yield strong MR (Methods); the negligible MR observed suggests that an additional, much more effective conducting channel is present. Second, the FS pocket deduced from the SdH oscillations would only account for 4–5% of the observed SdH data (Methods). Therefore, a separate FS of more conventional heavy fermions is required to account for the charge transport properties and large \( \gamma \) of the KM state. The heavy effective mass of these fermions (Methods) and the relatively high \( T > 0.5 \text{K} \) used for our measure-
ments would prevent the observation of their quantum oscillations, even in a very clean system (a combination of dilution refrigerator and pulsed magnet is under long-term development to overcome such challenges). The ‘bad-metal’ properties of the KM state, indicative of low charge carrier mobilities and thus create additional impediments towards the observation of quantum oscillations, give further support to the idea that the SdH effect is due to unconventional quasiparticles that are distinct from heavy electrons.

Our results imply an intriguing two-fluid picture in YbB$_{12}$ that includes (1) an FS pocket of quasiparticles obeying Fermi–Dirac statistics but contributing little to the transport of charge and (2) more conventional charged fermions. For brevity, we refer to the particles in (1) as neutral fermions (NFs). These NFs cause quantum oscillations in both KI and KM states, whereas the particles in (2) dominate the electrical transport properties and the low-$T$ heat capacity in the KM state. Under this two-fluid description, the I–M transition produces a sudden increase in the density of particles in (2), which changes from a thermally excited low-density electron gas in the KI state to a dense liquid of heavy, charged quasiparticles in the KM state. As a result, when $H$ increases past $H_{kI}$, the much enhanced number of states available close to the Fermi energy acts as a ‘reservoir’ into which quasiparticles from the NF FS can scatter as a ‘reservoir’ into which quasiparticles from the NF FS can scatter (for analogous situations in other materials, see refs. 31,32).

In this sense, the falling SdH frequency parameterized by equation (3) can be interpreted as an indication that the NF FS becomes less energetically favourable in the KM state; the availability of the KM ‘reservoir’ means that quasiparticles can transfer out and thus be depopulated. As the NF FS shrinks, the effective mass increases (Fig. 2c), possibly due to non-parabolicity of the corresponding band and/or field-induced modification of the bandwidth or inter-actions contributing to the effective mass renormalization$^{33}$. Also, the fact that the quantum oscillations caused by NF FS are strongly affected by the I–M transition (that is, the frequency becomes field-dependent on $H$) suggests that the quantum oscillations caused by NF FS are strongly affected by the I–M transition (that is, the frequency becomes field-dependent on $H$) suggests that the quantum oscillations caused by NF FS are strongly affected by the I–M transition (that is, the frequency becomes field-dependent on $H$) suggests that the quantum oscillations caused by NF FS are strongly affected by the I–M transition (that is, the frequency becomes field-dependent on $H$) suggests that the quantum oscillations caused by NF FS are strongly affected by the I–M transition (that is, the frequency becomes field-dependent on $H$).
against T (T*) from 4 K to 9 K. The inset illustrates the magnetic-field pulse and current pulse in our measurement in the time domain. ρb of represent the noise level for each MR curve.

...the quantum Hall effect. Whether this similarity is coincidental, however, is still an open question. In three-dimensional materials, an emergent gauge field, which can be felt by the charge carriers and manifests itself in transverse transport measurements, might originate from chiral spin textures in a skyrmion lattice phase. Nevertheless, such a topologically non-trivial magnetic structure has not yet been proposed for YbB12, and the potential physical importance of H* awaits further investigation.

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Fig. 3 | Temperature dependence of resistivity in the metallic state. a, The resistivity of the YbB12 sample at μH = 55 T plotted as a function of T. Both current I and magnetic field were applied along the [100] direction. The solid symbols are data measured using the pulsed-current technique (Methods). The solid symbols are data taken with a constant excitation in a 3He gas environment. A maximum in ρ at T* = 14 K. The dashed line is a linear fit of ρ(T) from 4 K to 9 K. The inset illustrates the magnetic-field pulse and current pulse in our measurement in the time domain. b, The solid line is a linear fit to the ρ(T) curve. The dotted line is a linear fit to the H* liquid data, showing the behaviour of the T* dependence below Tc = 2.2 K. The error bars in a and b represent the noise level for each MR curve. A1 and A2 are the T-linear coefficient and T2 coefficient of ρ (resistivity), respectively. c, The deviation of YbB12 from the KW relation. We use the value of the Sommerfeld coefficient γ reported for YbB12 in ref. 1. The data points for transition metals, d-electron oxides, and Ce- and U-based heavy fermions are taken from refs. 79,80, whereas the data for Yb-based compounds are taken from refs. 79,80."
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Methods

Sample preparation and pulsed-field facilities. YbB$_2$ single crystals were grown by the travelling-solvent floating-zone method$^{13}$. The two samples studied in this work were cut from the same ingot and shown to have almost identical physical properties, including the SDH oscillations below the I–M transition, in previous investigations$^{14}$. The YbB$_2$ sample characterized in the magnetostriction and magnetization (M) measurements corresponds to sample N1 in ref.$^1$ and crystal 2 in ref.$^1$, whereas the high-field MR was measured in the YbB$_2$ sample N3 in ref.$^1$, which is also crystal 1 in ref.$^1$. Both samples were used in the PDO experiments. The PDO data from sample N3 taken at a fixed field direction were published elsewhere$^{15}$. X-ray diffraction measurements on sample N1 show clear (006) series Bragg peaks (Extended Data Fig. 4) with no indication of crystalline impurity phases.

Magnetostriiction and M of YbB$_2$ samples were measured in a capacitor-driven 65 T pulsed magnet at the National High Magnetic Field Laboratory (NHMFL), Los Alamos. In the PDO and MR measurements, fields were provided by 65 T pulsed magnets and 75 T duplex magnet. Temperatures down to 500 mK were obtained using a He immersion cryostat. A He cryostat was also used for MR measurements above 1.4 K.

Magnetostriiction measurements. The linear magnetostriiction $\Delta L/L$ of YbB$_2$ was measured using fibre Bragg grating dilatometry$^{26}$. In our setup (Extended Data Fig. 1), the dilatometer is a 2-mm-long Bragg grating contained in a 125 µm telecom-type optical fibre. The oriented YbB$_2$ single crystal was attached to the section of fibre with the Bragg grating using a cyanocrylate adhesive. The crystalline axis was aligned with the fibre, which is also parallel to $H$. Thus, we measure the longitudinal magnetostriiction along the $a$ axis of cubic YbB$_2$. The magnetostriiction $\Delta L/L$ was extracted from the shift of the Bragg wavelength in the reflection spectrum$^{27}$. The signal from an identical Bragg grating on the same fibre with no sample attached was subtracted as the background.

In a paramagnetic metal, the high-field longitudinal magnetostriiction contains both $M^2$ and $M^3$ terms$^{14}$. In this sense, the power-law $H$ dependence of $\Delta L/L$ with an exponent of $-3.5$ (Extended Data Fig. 1b) is consistent with the weak superlinear $M$ in YbB$_2$ at 40 K (ref.$^{14}$). As $T$ decreases, $\Delta L/L$ decreases and a non-monotonic field dependence develops at 30 K (Extended Data Fig. 1b). We note that the fast suppression of $\Delta L/L$ coincides with the sharpening of the I–M transition in the derived susceptibility below 30 K (ref.$^{1}$), suggesting an additional energy scale in YbB$_2$ that is much lower than the Kondo temperature $T_K$.

$\Delta L/L_{\text{total}} = \Delta L/L_{\text{due to quantum oscillations}} + \Delta L/L_{\text{due to quantum magnetostriction}} + \Delta L/L_{\text{due to crystal anharmonicity}}$.

Magnetostriiction measurements. $M$ was measured using a compensated-coil susceptometer$^{48}$.$^{49}$. The 1.5-mm-bore, 1.5-mm-long, 1,500 turn coil was made of 50 gauge high-purity copper wire. The sample was inserted into a 1.3-mm-diameter non-magnetic ampoule that can be moved in and out of the coil. When pulsed fields are applied, the coil picks up a voltage signal $V = \delta (\delta M/\delta t)$, where $I$ is time. Numerical integration is used to obtain $M$ and a signal from the empty coil measured under identical conditions is subtracted. Pulsed-field $M$ data are calibrated using data of a Yb$_2$ sample of known mass measured in a vibrating sample magnetometer (Quantum Design) which is the same coil used in the experiments.

A metamatamagnetic transition occurs at $46.3 \pm 0.2$ T, coinciding with the onset of the step-like feature in the magnetostriiction. This observation further confirms the location of $H_{\text{I-M}}$ in our YbB$_2$ samples. At the highest $H$ used in this experiment, $M = 1.4\mu_B/\text{Yb}$, so that $M$ contributes only $-0.2\%$ of $B$. That is, we can ignore the $M$ term and equate $B$ to the external magnetic field, which is $B \approx \mu_B H$.

Radio-frequency measurements of resistivity using the PDO technique. The PDO circuit$^{12}$ permits convenient contactless measurements of the resistivity of metallic samples in pulsed magnetic fields. In our experiments, a 6–8 turn coil made from 40 gauge high-purity copper wire is tightly wound around the YbB$_2$ single crystal and secured using GE varnish. The coil is connected to the PDO, forming a driven LC tank circuit with a resonant frequency of 22–30 MHz at cryogenic $T$ and $H=0$. The output signal is fed to a two-stage mixer/filter heterodyne detection system$^{13}$, with mixer intermediate frequencies provided by a dual-channel BrK Precision function/arbitrary waveform generator. The intermediate frequency of the second mixer was 8 MHz, whereas the intermediate frequency of the first mixer was adjusted to bring the final frequency down to $\sim 2$ MHz. The resulting signal was digitized using a National Instruments PXI-5105 digitizer.

Considering all the contributions, the shift in PDO frequency $F$ due to $H$ is written as$^{12}$

$$\Delta F = -a \Delta I - \Delta R,$$

where $a$ and $b$ are positive constants determined by the frequency plus the capacitances, resistances and inductances in the circuit; $L$ is the coil inductance; and $R$ is the resistance of the coil wire and cables. In the case of a metallic sample, coil inductance $L$ depends on skin depth $\lambda$ of the sample. If we assume that sample magnetic permeability $\mu$ and coil length stay unchanged during a field pulse, we have $\Delta L \approx \pi \lambda/2$, where $\pi \lambda/2$ is the sample radius. At angular frequency $\omega$, the skin depth is proportional to the square root of resistivity $\rho$: $\lambda = \sqrt{2 \rho/\sigma}$.

Therefore, for a metallic sample, resonance shift $\Delta F$ reflects the sample MR and the detected quantum oscillations are due to the SDH effect. In YbB$_2$, the PDO measurement only detects the signal from the sample in the high-field $K$M state, that is, when $H > H_{\text{I-M}}$ (ref.$^1$). In the low-field $K$M state, the sample is so resistive that skin depth $\lambda$ is larger than sample radius $r$. As a result, $\Delta F$ mainly comes from the MR of the copper coil$^{12}$. The ‘dip’ in PDO is in fact a consistent with the sample radius and provides an alternative means to find $H_{\text{I-M}}$. We note that the $H_{\text{I-M}}$ value assigned to the onset of the ‘dip’ feature (Supplementary Table 1) is $\sim 0.2$ T lower than the metamagnetic transition shown in Extended Data Fig. 1.

Onsager relationship for a field-dependent FS. The Onsager relation$^{27}$ relates frequency $F$ of quantum oscillations to FS extremal area $A$: $F = \frac{4\pi}{\hbar} A$. Textbook derivations$^{12}$ invoke the correspondence principle to give an orbit-area quantization condition $(N + \lambda) \frac{2\pi}{\hbar} = A$, where $N$ is a quantum number and $\lambda$ is a phase factor. The derivation makes no assumptions about $A$ being constant, so for a field-dependent $A = A(B)$, we can write

$$\left( N + \lambda \right) \frac{2\pi e B_N}{\hbar} = A(B_N);$$

where $B_N$ is the magnetic induction at which the $N$th oscillation feature (peak, valley, etc.) occurs. Evaluating equation (6) for $N = 1$ and taking the difference gives

$$\frac{A(B_{N+1})}{B_{N+1}} - \frac{A(B_N)}{B_N} = \frac{2\pi e}{\hbar}.$$ 

From an experimental standpoint, in materials where $F$ is constant, a particular feature of a quantum oscillation is observed whenever $N + \lambda' = \frac{2\pi}{\hbar} B$, where $\lambda'$ is another phase factor. Allowing $F$ to vary ($F = F(B)$), evaluating the expression for $N = 1$ and taking the difference yields

$$\frac{F(B_{N+1})}{B_{N+1}} - \frac{F(B_N)}{B_N} = 1.$$ 

A comparison of equations (7) and (8) shows that the Onsager relation still holds, that is, $F(B) = \frac{2\pi e}{\hbar} A(B)$.

In Fig. 2a, $B^* = \mu_0 H_{p}$ is 41.6 T is subtracted from the applied field to yield linear Landau-level index diagrams for $\theta \leq 20^\circ$. The resulting fit is described by $N + \lambda = \frac{2\pi e B}{\hbar} = \frac{2\pi}{\hbar} (F/B)$ (equation (3)), which can be written in terms of a B-dependent frequency $N + \lambda = \frac{2\pi e}{\hbar} B$, where $F(B) = \frac{2\pi e}{\hbar} B$ (equation (3)). Following the reasoning given above, this is associated with a $B$-dependent extremal area

$$A(B) = \frac{A_0}{B - B_0},$$

where $A_0 = \frac{2\pi e}{\hbar} B_0$.

Note that we have not considered the Zeeman splitting of the peaks; therefore, equations (3) and (9) describe the average of the spin-up and spin-down components.

Spin-splitting parameter. The spin-splitting parameter $\lambda$ is introduced to describe the Zeeman splitting of the quantum oscillation peaks/valleys. In the most straightforward scenario, the magnetic inductions $H_{\text{I-M}}(\lambda)$ at which the spin-down (spin-up) component of the $N$th Landau tube reaches the FS are given by

$$\frac{F}{B_N} = N + \lambda \pm \frac{1}{2} \frac{\pi e}{\hbar}.$$ 

Therefore, $\lambda$ causes a term in addition to the phase factor $\lambda$, which has opposite signs for the spin-down and spin-up Landau sheets. According to equation (10),
the value of $S$ can be obtained by multiplying the interval of the spin-split peak/valleys (in 1/2B) by the oscillation frequency $F$. In Fig. 2d, $S$ for nonlinear (Fig. 1c) and linear (Fig. 2a) Landau-level plots are calculated as $S = F_{1}(1/B_{j}^{*} - 1/B_{i}^{*})$ and $S = F_{2}(1/B_{j}^{*} - B_{i}^{*})$ and $1/(B_{j}^{*} - B_{i}^{*})$, respectively. Here $F_{1,2}$ represent the oscillation frequency directly extracted from the slope of the nonlinear Landau-level plot (Supplementary Information). We mention that these two approaches yield the same values of $S$ (Fig. 2d; data for $\theta = 10.9^\circ$).

As pointed out by Shoenberg, the effect of Zeeman splitting on oscillation extrema is determined by the cyclotron mass $m^*$ and g factor:

$$S = \frac{1}{2} \frac{m^*}{m_0}$$

which provides a simple method to evaluate the g factor in the system. As shown in Fig. 2d, with a g factor of 0.084, the field dependence of $S$ can be adequately mapped onto the mass enhancement. Such behaviour does not only prove the validity of our Zeeman splitting analysis of the SdH effect but also reveals an unusually small g factor that may reflect the exotic nature of the corresponding quasiparticle band.

Resistivity measurements in the KM state. From $\mu B = 0$ to 60 T, the resistivity of YbB$_2$ decreases by five orders of magnitude. As a result, MR measurements in pulsed fields are challenging. If a constant current is used during the entire field pulse, even the signal-to-noise ratio is poor in the KM state or the high resistance in the KI state causes serious Joule heating. To resolve this issue, we developed a pulsed-current technique. The experimental setup is shown in Extended Data Fig. 3a. A ZFSWHFA-10 isolation switch is used to apply current pulses with widths less than 10 ms. The switch is controlled by two square-wave pulses generated by a B&K Precision model 4665 dual-channel function/arbitrary waveform generator triggered by the magnet pulse. Thus, a relatively large electric current (2-3 mA) can pass through the sample during a narrow time window within the field pulse (Fig. 3a, inset).

Current is applied only when YbB$_2$ enters the KM state. In our low-T MR measurements, the switch turns on at 47 T during the upsweps and turns off at 47 T during the downsweps. To reduce heating due to eddy currents, we measured the longitudinal MR of a needle-shaped sample (length of 6.5 mm between voltage leads; cross-sectional area, 0.33 mm$^2$). As shown in Extended Data Fig. 3b, the downsweps still suffer from some Joule heating. Hence, in the main text, we focus on the upsweps, which show very weak longitudinal MR and should reflect the intrinsic electrical transport properties of the KM state of YbB$_2$. By following the experiments, we found that when we attempted to stabilize the target temperature by applying a high heater excitation in an environment of dense He vapour, the temperature reading according to the thermometer could be higher than the actual sample temperature. We removed all the data points measured using this procedure.

According to equation (5), $A(\Delta)$ for the pocket detected in the PDO measurement shrinks by ~45% from 50 T to 60 T. Assuming a spherical FS, this corresponds to an ~60% reduction in quasiparticle density $n$. Meanwhile, the cyclotron mass increases by ~60% (Fig. 2b, inset). Consequently, a textbook Drude expression $\rho = m^* n/e^2 (\tau$ is the relaxation time) predicts that the resistivity would increase by a factor of ~4 from 50 T to 60 T if the electrical transport is dominated by this pocket. In sharp contrast, almost no MR is observed in this field range (Extended Data Fig. 3), indicating the negligible contribution of this FS pocket to the electrical transport.

For the Sommerfeld coefficient

$$\gamma = \frac{\pi^2 k_B^2 m^* k_B}{3 e^2}$$

(12)

where $k_B$ is the Fermi vector. As for the FS pocket detected in the SdH measurement, equation (3) gives $F(35.5 \Omega T) = 231.9 T$. In a spherical FS model, $F = \hbar^2 k_B^2 / 2e$ and $m^* k_B^2 = 2(2\pi)^3/3(\pi n)$, the Sommerfeld coefficient can be written as

$$\gamma = \frac{\pi^2 k_B^2 m^* k_B}{3 e^2}$$

(13)

Taken together, equations (12) and (13) yield $k_B = 2.15 \text{nm}^{-1}$ (corresponding to $n \approx 3.36 \times 10^{20} \text{cm}^{-3}$) and a rather large effective mass of $m^* = 90 m_0$. Such a heavy mass is unusual for Yb-based mixed-valence compounds, but explains the anomalously high $\gamma$ as well as the absence of SdH oscillations due to these quasiparticles in the PDO response.

SdH oscillations due to charge-neutral quasiparticles. Magnetic quantum oscillations are observed in both $M$ and $\rho$ in the KM state of YbB$_2$, with their behaviours following the description of the LK formula. The size of $M$ and the effective mass of quasiparticles inferred from the oscillations are consistent with the fermion-like contribution to the thermal conductivity. A natural explanation—given the high electrical resistivity—is that the thermal conductivity and quantum oscillations are due to charge-neutral fermions. (We mention that a hypothetical model of the SdH effect in a system containing gapless neutral fermions and gapped charged bosons has been established, which predicts the LK behaviour above a certain $T$.)

Of the two oscillatory effects observed in the KM state, the dHvA effect is the more fundamental. As pointed out by Lifshitz, Landau and others, it involves oscillations in a thermodynamic function of state (M) that may be related to the electronic density of states with a minimum number of assumptions. The fact that $M$ in the KM state oscillates as a function of $H$ can only be due to the oscillation of the fermion density of states and consequently their free energy.

Even in conventional metals, quantum oscillations in $\rho$ are harder to model quantitatively. A starting point was suggested by Pippard’s theory of quasiparticles scatter depends on the density of states available via Fermi’s golden rule. Hence, if the quasiparticle density of states oscillates as a function of $H$, the scattering rate $\tau^{-1}$ and consequently $\rho$ will also proportionally oscillate, leading to the SdH effect. Before modifying this idea to tackle the SdH effect in the KM state, we remark that the $H$-dependent frequency of oscillations in the KM state suggests that the exchange of fermions between charge-neutral and conventional FS sections occurs readily, probably via low-energy scattering. This is also supported by the $T$-linear resistivity in the KM state (Fig. 3a). The rate at which this scattering happens will obviously reflect the joint density of fermion states.

Returning to the KM state of YbB$_2$, $\rho$ is thought to be due to charge carriers thermally excited across the energy gap, plus contributions from states in the gap that lead to $\rho$ saturation at low $T$. Following the precedent of the KM state, it is likely that fermions in the KM state scatter back and forth between the charge-neutral states and the more conventional bands. The situation is more complicated than Pippard’s original concept because the scattering between two bands is involved (for example, ref. 19). Nevertheless, the amount of scattering will be determined by the joint density of states; because the density of states of the neutral quasiparticles oscillate with $H$, so does $\rho$. Unlike conventional metals, the density of charge carriers depends on $T$, so the amplitude of SdH oscillations in the KM state follows a different $T$ dependence from that of the dHvA oscillations, as observed in experiments.

As described by equation (4), the PDO in the KM state is determined by the skin depth, that is, the conductivity. As described above, the scattering of fermions occurs between the charge-neutral states and more conventional bands; in the KM state, our conductivity results and heat capacity data of others suggest that the latter is an FS of heavy charged fermions. The SdH effect in conductivity caused by the oscillatory density of states of the neutral quasiparticles in magnetic fields is detected in the PDO experiments. By contrast, any intrinsic quantum oscillations due to the metallic FS are suppressed by a combination of the very heavy effective mass and relatively high $T$ ($\approx 20 K$) of our measurements.

Data availability
The data that support the plots within this paper and other findings of this study are available from the corresponding authors upon reasonable request.

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Author contributions

F.I. grew the high-quality single-crystalline samples. Z.X., L.C., K.-W.C., C.T., Y.S., T.A., H.L., F.B., J.S. and L.L. performed the pulsed-field PDO and resistivity measurements. Z.X., T.A. and J.S. performed the pulsed-field magnetometry measurements. L.C., C.T. and M.J. performed the pulsed-field magnetostriction measurements. Z.X., K.-W.C., Y.K., Y.M., J.S. and L.L. analysed the data. Z.X., Y.M., J.S. and L.L. prepared the manuscript.

Competing interests

The authors declare no competing interests.

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Extended Data Fig. 1 | The magnetostriction and magnetization of YbB₁₂.  

**a.** The experimental arrangements for magnetostriction measurements. An optical fiber containing fiber Bragg gratings serves as the strain gauge. The single-crystal sample of YbB₁₂ with dimensions 2.3×1.0×0.2 mm³ is glued to the fiber Bragg grating section of the fiber. A 60 T pulsed magnetic field is applied along the direction of the optical fiber. The longest edge of the single crystal is parallel to the crystallographic a axis, therefore in this arrangement the parallel strain $\Delta L / L \parallel [100]$ is measured.  

**b.** Field dependence of the linear magnetostriction $\Delta L / L$ along the a axis at different temperatures ranging from 10 K to 40 K. The inset shows the high-temperature magnetostriction curves in a logarithmic scale. The dashed line is a power-law field dependence with an exponent $n = 3.5$.  

**c.** The linear magnetostriction measured at temperatures below 5 K.  

**d.** Magnetization of the same YbB₁₂ single crystal measured at $T = 0.64$ K using a compensated-coil susceptometer (see Methods). Field was applied along the [100] direction. The gray dashed line indicates the consistency of the step-like decrease in the linear magnetostriction in c and the abrupt increase of magnetization in d at the field-induced I-M transition, $\mu_0 H_{I-M} = 46.3$ T.
Extended Data Fig. 2 | The angle, temperature and field dependence of the SdH effect. a, Oscillatory component of the PDO frequency, $\Delta f$, obtained after a 4th-order polynomial background subtraction from raw data at different tilt angles. The SdH effect is weaker in the field downsweeps (short-dashed lines) than in the upsweeps (solid lines), probably due to sample heating. b, $\Delta f$ taken at $\theta=10.9^\circ$ at different temperatures. The field-dependent cyclotron energies shown in Fig. 1d are obtained by fitting the frequency difference between adjacent peaks and valleys to the LK formula. Horizontal bars denote the peaks and valleys used for the fits in Fig. 1d. The inset displays the interval between spin-up and spin-down SdH peaks in inverse magnetic field, as a function of $\mu_0 H$. If the effective mass and the g-factor are field independent, this value is expected be a constant (see Methods). c, PDO frequency measured up to 75 T in the Duplex Magnet; only the lowest pair of Landau levels, that is, $N=2\pm1$, are shown. The inset shows the small amplitude differences of $N=2\pm1$ levels between $T=0.57$ K and 1.51 K, suggesting a nondiverging quasiparticle mass up to at least 72 T. d, MR measured by the pulsed current technique (see Methods) in the Duplex magnet. A downward kink is observed at 68 T and coincides with the $N=2\pm1$ sublevel. This slope change may imply a crossover to an unknown high-field state in YbB$_6$. 
Extended Data Fig. 3 | The pulsed current technique. a, A schematic diagram of the pulsed current technique used for measuring the MR of YbB$_2$ samples in the KM state (see Methods). The current pulse is triggered with a precisely controlled time delay after the start of the magnetic field pulse. In the measurements we adjust the value of $R_s$ to maintain the excitation amplitude 2-3 mA at different temperatures. b, The longitudinal MR of YbB$_2$ in the KM state below 10 K, measured with both field and current applied along the crystallographic [100] direction. The data taken at 1.44 K, 2.00 K and 3.09 K are measured with the sample in 4He liquid, whereas other curves are measured in 3He liquid or vapour. The vertical dashed line marks the 55 T field at which we extract resistivity values from the upsweeps for the temperature dependence analysis (see Methods).
Extended Data Fig. 4 | X-Ray diffraction. X-ray diffraction pattern of the YbB₁₂ single crystal measured with Cu Kα radiation. Only the (00l) series Bragg peaks are observed. The lattice parameter is calculated to be \( a = 7.47 \) Å.