

Magnetoelastoresistivity in FeSe

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FeSe stands out among iron-based superconductors due to its extended nematic phase without the onset of long-range magnetic order. While strain-dependent electrical resistivity has been extensively explored to probe nematicity, its influence on magnetotransport properties remains less understood. In this work, we present measurements of the magnetoelastoresistivity in FeSe as a function of temperature and applied magnetic field. Using a minimal multiband Boltzmann model for transport we derive analytical expressions that capture the magnetic behavior of the whole set of experimental data both in the paramagnetic and in the nematic phase. These findings indicate that a multiband framework can robustly describe the magnetoelasto-transport properties in FeSe and arguably in other iron-based superconductors.

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I. INTRODUCTION

Since their discovery, iron-based superconductors (Fe-SCs) have provided a rich platform to explore the interplay between structural, magnetic, and orbital degrees of freedom in unconventional superconductivity [1]. Among the various emergent phenomena, the so-called nematic phase, defined by the spontaneous breaking of the crystal rotational symmetry while preserving translational invariance, has attracted particular attention due to its close connection with magnetism and superconductivity [2–4]. In Fe-SCs, nematicity typically emerges before or alongside magnetic ordering, suggesting a close interplay between the two. Nematic fluctuations can enhance magnetic interactions, while magnetic fluctuations may contribute to the onset of nematicity and their mutual influence is believed to play a role in shaping unconventional superconductivity [2–4]. In this context, the FeSe compound stands out as a notable exception. It exhibits an extended nematic phase below $T_S \approx 90$ K without developing long-range magnetic order, and hosts bulk superconductivity below a critical temperature $T_C \approx 9$ K [5,6]. Although the absence of long-range magnetism does not preclude the presence of

magnetic fluctuations, FeSe offers a unique opportunity to investigate the relationship between nematicity and superconductivity in the absence of magnetic order. There is nowadays a general agreement on the electronic nature of the nematic state, however, the exact origin of nematicity in FeSe is still under debate. The lack of magnetic order motivated a scenario based on orbital or charge ordering [7–9], while the detection of sizable spin fluctuations [10–13] triggered an intense investigation on the interplay between spin and orbital degrees of freedom [12,14–18], leading to an alternative picture of nematicity driven by orbital-selective spin fluctuation [6,19,20].

A powerful experimental technique to probe the coupling between nematic and electronic degrees of freedom is the measurement of the strain (ϵ) dependent electrical resistivity (ρ), or elastoresistivity, $ER = \frac{d\rho}{d\epsilon}$ with $\eta = \frac{\rho(\epsilon) - \rho(\epsilon=0)}{\rho(\epsilon=0)}$, which has been widely used over the past decade [21–28]. In particular, the measurement of the ER as a function of temperature and doping allowed to identify the electronic origin of nematicity in Fe-SCs and to reveal a correlation between the enhanced nematic fluctuations and the increase of the superconducting critical temperature at the optimal doping. Curiously, much less attention has been devoted to the study of the magnetoelastoresistivity (MER), the measurement of ER under an applied magnetic field, which can offer complementary insights into the symmetry-breaking mechanisms by probing the magnetic field dependence of strain-induced transport anisotropy. To fill this gap, we present a detailed investigation of MER in FeSe as a function of temperature and applied magnetic field. We analyze the experimental data using the Boltzmann transport theory, taking into account the multiband nature of the electronic structure

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of FeSe and the transport anisotropy induced by nematicity below T_S . While the effects of electronic correlations and bands' orbital composition are not explicitly modeled, they are effectively incorporated at the microscopic level in the x/y band conductivities, used in defining the fitting parameters. Our analysis captures the essential behavior observed in the magnetotransport experiments. Notably, two key features of the MER data emerge directly from the analytical expressions, without any fine-tuning of fitting parameters: (i) the field-independent behavior above T_S for out-of-plane (i.e., parallel to the crystallographic c -axis) fields, and (ii) the absence of magnetic-field dependence for in-plane (i.e., perpendicular to the crystallographic c -axis) fields. These results demonstrate that a minimal two-band model robustly captures the essential magnetoelastotransport features in FeSe, suggesting a unified framework for analyzing MER in other Fe-SCs.

II. EXPERIMENTAL METHODS

FeSe single crystals were synthesized according to the procedure reported in Ref. [29]. The samples used for the experiments are two specimens cleaved from the same single crystal with an approximate size of $1.5 \times 1.5 \times 0.03 \text{ mm}^3$. All the measurements were performed using homemade probes inserted in an Oxford Instruments cryostat, equipped with a 15/17 T magnet. The transport characterization has been carried out adopting a standard 4-wire configuration, in which the electrodes consist of 50 μm -thick silver wires attached to the sample with silver paints. The MER measurement setup was described in Refs. [24,27]. The sample was glued on one side a commercial piezoelectric stack actuator (by Piezomechanik GmbH) endowed with a low-temperature coating to perform measurements in cryogenic environment. On the other side of the piezostack a strain-gauge was glued to monitor the actual strain in the experiments. All the measurements were performed in a limit of small strain ($\epsilon < 5 \times 10^{-4}$), characterized by a linear dependence of the resistivity on strain and by the complete reversibility of the response. The stress was applied to the sample by aligning the $[1\ 1\ 0]$ crystalline direction of FeSe samples to the piezomain axis of motion [Fig. 3(a)]. Additional measurements with the piezoaligned along the $[1\ 0\ 0]$ crystalline direction of a FeSe sample are shown in the Supplementary Material (SM) [30]. With this procedure the measured ER is a combination of the B_{2g} mode, associated to the nematic phenomenology, and the A_{1g} mode, whose relative contribution to the elastoresistance is negligible in FeSe [31]. The detailed measurement protocol is described in the SM [30]. The magnetic field in MER measurements has been oriented both out of plane and in plane (i.e., parallel and perpendicular to the crystallographic c -axis).

III. RESULTS

The temperature-dependent resistivity (normalized to the value at $T = 280 \text{ K}$) $\rho/\rho_{280 \text{ K}}$ of two monocrystalline FeSe samples (S1 and S2) from the same batch is displayed in Fig. 1. The two curves almost overlap and exhibit the characteristic signatures of the structural and superconducting transitions at $T_S = 89 \text{ K}$ and $T_C = 9 \text{ K}$, respectively, in agreement with previous reports [5,32]. Figure 2 shows the

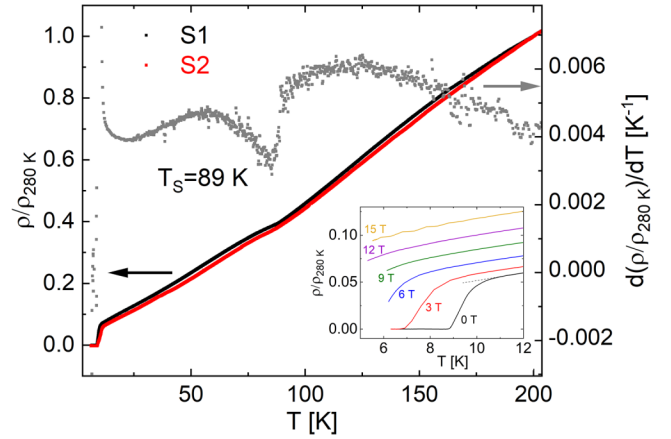


FIG. 1. Temperature dependence of the resistivity of samples S1 and S2 normalized to the values at 280 K. The temperature-derivative is shown in gray and clearly identifies T_S and T_C . Inset: Temperature dependence of the resistivity of sample S1 across the superconducting transition at several magnetic fields.

results of magnetotransport measurements of sample S1. In particular, Fig. 2(a) presents the magnetoresistivity (MR) $\frac{\Delta\rho(B)}{\rho(0)} \equiv \frac{\rho(B) - \rho(0)}{\rho(0)}$ as a function of the magnetic field B , applied along the crystallographic c -axis, at several temperatures. The MR is relatively small in the tetragonal phase and it grows below T_S , reaching the remarkable value of 90% at 15 K and 15 T, without signatures of saturation.

The field dependence of the Hall resistivity ρ^{xy} is shown in Fig. 2(b). Above T_S , ρ^{xy} is linear-in-field and negative, while at $T = 90 \text{ K}$ ρ^{xy} turns positive with a larger slope. With further decreasing the temperature, in the nematic phase, ρ^{xy} becomes strongly nonlinear, with field-induced sign changes. The inset shows the temperature dependence of the Hall coefficient R_0^H , thereby evaluated as $R_0^H = \frac{d\rho^{xy}}{dB}|_{B \rightarrow 0}$, which exhibits a first sign change at $T \approx 100 \text{ K}$ and a second sign change at $T \approx 60 \text{ K}$, in agreement with previous reports [33,34].

Figure 3(c) presents the MER of S1 as a function of the temperature for different magnetic fields applied along the crystallographic c -axis. The zero-field ER is consistent with existing literature [25,35,36]. Above T_S , ER is positive and exhibits the characteristic Curie-Weiss-like behavior, generally interpreted as the contribution of electron-driven nematic fluctuations. The extracted Curie-Weiss temperature is $T^* \approx 66 \pm 6 \text{ K}$ [see inset of Fig. 3(c)]. Below T_S , the ER drops and becomes negative for $T < 50 \text{ K}$.

Within our experimental resolution, no substantial field-induced shift of T_S is observed and the Curie-Weiss fit of the MER at 15 T returns the value of $T^* = 70 \pm 8 \text{ K}$, in good agreement with the one obtained from the fitting of the zero-field ER. Below T_S , the MER strongly increases towards positive values and it shows a pronounced upturn with decreasing T . We tested the consistency of the sign change through the isothermal measurement of MER versus B at $T = 20 \text{ K}$, shown in Fig. 3(b): the magnetic-field dependence is substantially linear for $B > 4 \text{ T}$, with a departure from linearity only at lower fields.

Figure 3(d) shows the MER of S2 with $B = 0$ and 15 T, applied both parallel and perpendicular to the crystallographic

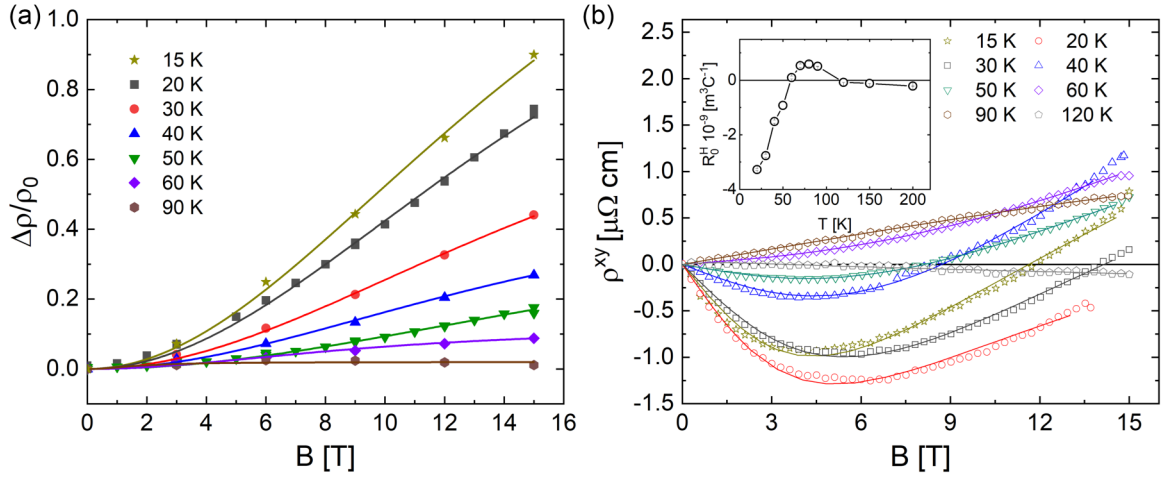


FIG. 2. (a) MR of sample S1 as a function of the applied magnetic field B , for different temperatures. (b) Hall resistivity ρ^{xy} versus B of samples S1, measured at different temperatures. Inset: Hall coefficient R_0^H , here derived as the slope of the linear fits in the low- B limit. In both panels (a) and (b), the theoretical fit using Eqs. (3) and (2) are shown as solid lines on top of the experimental data (symbols). Fitting parameters are reported in Appendix B.

c -axis of the sample. Both the ER and the MER at 15 T, parallel to the c -axis, are in good agreement with the measurement on sample S1. In addition, in this second set of data we verified that a magnetic field oriented in plane has no effect on the MER even below T_S , with the MER curves at ± 15 T overlapping over the zero-field ER.

IV. DISCUSSION

Our goal is to derive a minimal theoretical framework that consistently describes the full set of magnetoelastoresistivity measurements presented for FeSe. A natural starting point for modeling electron transport is the Boltzmann equation, which provides a semiclassical framework for charge conduction. In Fe-SCs, this approach must be extended to a multiband formulation to account for the complex electronic structure that governs their transport properties. ARPES experiments on FeSe have established the presence of multiple hole-like and electron-like pockets and have shown that the Fermi surface undergoes a pronounced reconstruction across the structural transition at T_S , accompanied by a marked in-plane anisotropy of the low-energy electronic structure [4,19,35,37–41]. These features make a single-band description inadequate for transport and motivate a multiband model with anisotropic band conductivities below T_S . To keep the discussion as simple and transparent as possible, we therefore consider a two-band model and account for the onset of nematic order below T_S by introducing an x/y anisotropy in the band conductivities. Importantly, we do not make any assumptions about the microscopic origin of nematicity. To preserve generality, we work in the band basis without explicitly introducing orbital degrees of freedom. As we will show, this minimal model is sufficient to reproduce the key features of the MER data without invoking orbital-selective effects. While the orbital composition of the bands is undoubtedly important at the microscopic level, particularly in determining the anisotropic conductivities in the nematic phase [42,43], here these effects are effectively absorbed into the fitting parameters that define the x/y band conductivities in our model.

In what follows, we summarize the main steps of the derivation and report only the equations relevant for the discussion. A detailed derivation of the Boltzmann transport theory for a two-band system in both the paramagnetic and nematic phases is provided in Appendix A.

Assuming that the electronic scattering rate is independent of the magnetic field, the resistivity tensor in the presence of an out-of-plane magnetic field (denoted as B_z) takes the form

$$\hat{\rho} = \begin{pmatrix} \rho^x & B_z R^H \\ -B_z R^H & \rho^y \end{pmatrix}, \quad (1)$$

where $\rho^{x/y}$ and $\rho^{xy} = R^H B_z$ are expressed in terms of the components of the band resistivity tensors $\hat{\rho}_n$ ($n = 1, 2$ is the band index). The explicit expressions for $\rho^{x/y}$ and R^H in the two-band model are given in Eqs. (A9)–(A11). The Hall coefficient is given by

$$R^H = \frac{a_0 + a_1 B_z^2}{1 + a_2 B_z^2}, \quad (2)$$

where a_0 , a_1 , and a_2 are functions of the components of $\hat{\rho}_n$. This expression holds in both isotropic and anisotropic cases, only the explicit forms of a_0 , a_1 , and a_2 differ, as detailed in Appendix A. Similarly, the MR can be written as

$$\frac{\Delta\rho(B)}{\rho(0)} = \frac{1 + b_0 B_z^2}{1 + b_1 B_z^2} - 1, \quad (3)$$

where b_0 and b_1 are also functions of the band resistivity tensors, with their explicit forms in the isotropic and anisotropic cases provided in Appendix A.

To analyze the MR data, we treat $a_{0,1,2}$ and $b_{0,1}$ as fitting parameters. The results of the fitting procedure are shown in Figs. 2(a) and 2(b), alongside the experimental data. The corresponding parameter values are listed in Appendix B. A good fit is obtained using a physically reasonable set of parameters. It is worth noting that the fit assigns different values to b_1 and a_2 , although within our model these parameters are defined by the same expression (see Appendix A). We attempted a simultaneous fit of MR and ρ^{xy} under the

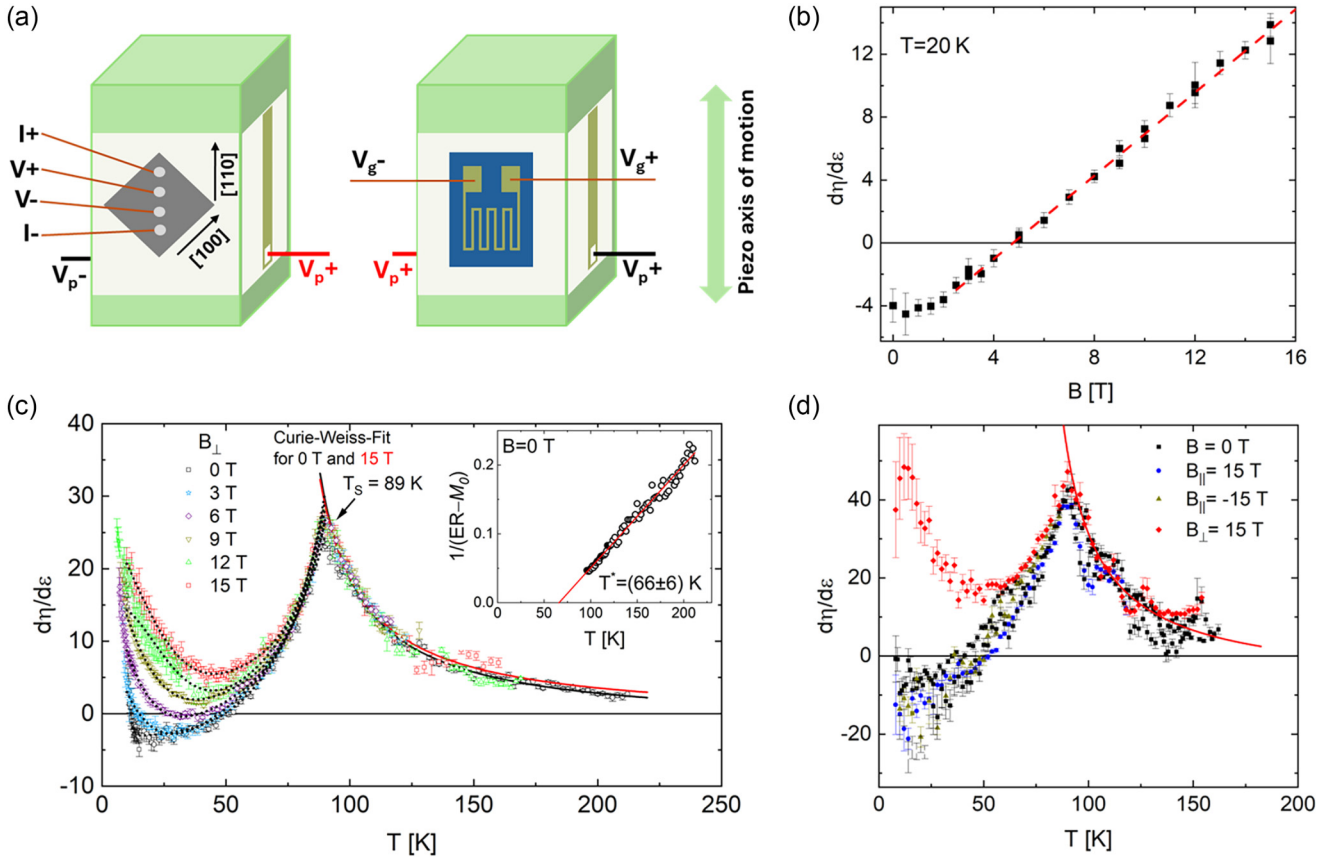


FIG. 3. (a) Schematic representation of the sample mounting on a piezoelectric device. The sample is glued on the front side on the piezostack, with the crystallographic $[1\ 1\ 0]$ direction oriented along the axis of motion. A strain gauge is mounted on the back side of the piezostack. (b) Magnetic-field dependence of the ER of sample S1 at $T = 20$ K, these are the type of data we fitted using Eq. (4). (c) ER on sample S1 in out-of-plane fields, from 0 T (black) to 15 T (red), in 3 T steps, shows a systematic increase with increasing field in the orthorhombic phase. On top of the experimental data (symbols) at $T < T_s$, we show the curve obtained by evaluating Eq. (4) at the corresponding field values using the temperature-dependent fitting parameters provided in Appendix B (dashed lines). The behavior in the tetragonal phase is instead reproduced by the Curie-Weiss-like function $ER = M_0 + M_1/(T - T^*)$. Black and red solid curves interpolate the data set at 0 T and 15 T, respectively. Inset: Linear fitting of the $1/(ER - M_0)$ versus T curve at 0 T to extract the value of $T^* = (66 \pm 6)$ K. (d) MER of sample S2 for different magnetic field configurations: zero-field (black), 15 T and -15 T in-plane (blue and brown) and 15 T out-of-plane (red). No effect of magnetic fields is visible for in-plane fields at any temperature.

constraint $b_1 = a_2$. As expected, the inclusion of an additional constraint slightly worsens the fit quality, particularly at low temperatures. However, the overall trend of the fitting parameters remains unchanged (see Appendix B), which supports the internal consistency of our minimal two-band model in capturing the essential features of the magnetoelastotransport response in FeSe. Importantly, since the parameters b_1 and a_2 primarily affect the MR and ρ^{xy} in the high-field regime, their discrepancy may point to a limitation of the standard perturbative treatment of the Boltzmann equation via the Jones-Zener expansion [44]. This approximation assumes the scattering rate to be independent of the magnetic field. A more rigorous solution of the Boltzmann equation, as discussed in Ref. [45], accounts for the magnetic field's influence on the time evolution of the wave vector, thereby modifying the scattering rate along the trajectory. In such a framework, b_1 and a_2 are no longer expected to be equal, however, to undertake this approach one would be forced to abandon the analytical treatment in favor of a fully numerical analysis.

The resistivity anisotropy is defined as $N = 2(\rho^x - \rho^y)/(\rho^x + \rho^y)$, where the specific forms of $\rho^{x/y}$ for the two-band model are given in Eqs. (A9) and (A10). It is easy to verify that in the limit of small strain $N = \eta$ [22]. To model the effect of strain ϵ on transport, we assume that $\rho_n^{x/y} = \rho_n^{x/y,0} \pm \frac{1}{2}\gamma_n\epsilon$, where γ_n is the strain-coupling coefficient of band $n = 1, 2$. Above T_s , symmetry requires the strain-free resistivity to be isotropic for each band, i.e., $\rho_n^{x,0} = \rho_n^{y,0} \equiv \rho_n^0$, so that the resistivity anisotropy is proportional to strain and symmetric with respect to x/y . In this case, the MER reduces to

$$\frac{\partial N}{\partial \epsilon} = c_0 \frac{1 + c_1 B_z^2}{1 + c_2 B_z^2}, \quad (4)$$

where c_0, c_1, c_2 are functions of the components of the band resistivity tensors $\hat{\rho}_n$, and of the strain coupling coefficients γ_n , whose explicit expressions are given in Appendix A. The analytical expression in Eq. (4) successfully captures the experimental behavior. In particular, assuming $\gamma_1 \sim \gamma_2$, we

find $c_1 \sim c_2$ and $c_0 \sim \gamma/\rho$, implying that the MER becomes independent of the magnetic field above T_S , consistent with the experimental observation [see Figs. 3(c) and 3(d) and the discussion below Eq. (A28)]. The general expression valid in the nematic regime is given in Eq. (A30) of Appendix A and depends on a total of six c_i coefficients. As discussed in Appendix B, we first analyze the magnetic-field dependence of the data at fixed temperature using the full expression and find that retaining all allowed coefficients does not significantly improve the agreement with the data. To avoid superfluous fitting parameters, we therefore restrict the analysis to the reduced parametrization of Eq. (4), involving only $c_{0,1,2}$. The result of the procedure is shown in Fig. 3(c), together with the experimental data, while the fitting parameters are provided in Appendix B. A good agreement with the experimental data is obtained using a set of parameters consistent with those used to analyze the MR data. The success of the reduced parametrization indicates that the overall strain-free resistivity is effectively fourfold symmetric at the level of the measured macroscopic response, even in the nematic phase, where the symmetry is locally broken within each domain. A natural interpretation of this result is that, in the absence of applied strain, the two nematic-domain orientations have approximately balanced populations, so that the measured resistivity represents an effectively C_4 -symmetric domain average. This interpretation is consistent with the small, linear, and reversible strain regime explored in our measurements in which the applied strain modifies the relative domain populations without producing a persistent domain imbalance once the strain is released.

So far, we consider the transport properties of the system in the presence of an out-of-plane magnetic field B_z , as described by Eqs. (1)–(4). In Fig. 3(c) we show the experimental results obtained under an in-plane magnetic field and find that, in this case, the MER is independent of the magnetic field at all temperatures. To theoretically analyze this situation we derive the Boltzmann equation in the presence of an in-plane magnetic field. The derivation follows the same steps as for the B_z case. However, if we assume that the system is nearly two-dimensional (with a larger effective mass along the c crystallographic axis), the resistivity tensor in Eq. (1) becomes nearly diagonal due to the negligible orbital-magnetic response (see Appendix A). As a consequence, the magnetic dependence of the MER, as expressed in Eq. (4), is governed by almost vanishing c_1 and c_2 coefficients. This leads to a MER that is independent of the in-plane magnetic field at all temperatures, in agreement with the experimental findings.

Summarizing the theoretical analysis, we show that the magnetoelastotransport measurements in FeSe can be simultaneously and consistently described using the Boltzmann equation within a minimal two-band model. The nematic order is incorporated phenomenologically by introducing an explicit x/y anisotropy in the band resistivities below T_S . Remarkably, while the orbital composition of the electronic bands is certainly important for the microscopic determination of $\rho_n^{x/y}$ in the nematic phase, it is not a crucial ingredient for achieving a consistent description within the present framework. Working with a minimal band model, without introducing additional degrees of freedom, allows for an analytical treatment of the relevant physical quantities and

provides clear physical insight. In particular, key features of the MER magnetic-field dependence emerge naturally from inspection of the analytical expressions, without requiring fine-tuning of the fitting parameters. Finally, using a limited number of parameters, we could successfully fit the full set of experimental data.

We emphasize that the present model focuses on describing the magnetic-field dependence of the magnetoelastotransport measurements, while the temperature dependence of the various quantities is treated phenomenologically. Specifically, we account for the x/y anisotropy of the band resistivities $\rho_n^{x/y}$ in the nematic phase, but to retain generality and preserve analytical tractability, we do not make any assumptions about the microscopic origin of nematicity. To theoretically describe the temperature dependence of MER near the nematic transition, including the expected $\gamma \sim 1/(T - T_S)$ behavior, one would need to define how electrons couple the nematic order parameter to the magnetic field within the model. However, the experimental behavior of MER around T_S , including the observation that T_S remains unchanged under magnetic field, suggests that this coupling is weak and can be neglected to first approximation.

The superconducting critical temperature T_C is the other relevant energy scale that should be incorporated to fully capture the temperature dependence of MER. In particular, the pronounced upturn of MER approaching T_C in Figs. 3(c) and 3(d) is consistent with an additional contribution from superconducting fluctuations to transport [46,47]. This interpretation is qualitatively supported by the temperature dependence of the resistivity (inset of Fig. 1), where a downturn in ρ/ρ_{280K} begins around $T \approx 11$ K, well above T_C . As expected, this fluctuation-induced contribution is suppressed with increasing magnetic field [48]. A detailed analysis of the regime near T_C will be the subject of future work.

As a final remark, we notice that, although in this work we did not address the temperature dependence of the experimental quantities, the fitting of MR, Hall resistivity, and MER at different temperatures provides valuable information on the temperature evolution of the fitting parameters and therefore can be used in future studies, e.g., to identify the relevant scattering mechanisms within a more microscopic framework.

V. CONCLUSION

In this work we simultaneously examined the transport properties in the magnetic field of the iron-based superconductor FeSe, across the nematic transition. Beside the conventional magnetotransport (Hall effect and magnetoresistance), we explored the magnetoelastoresistance, namely, the elastoresistance measured in magnetic field, which has been rarely experimentally addressed.

The overall behavior of the Hall and magnetoresistance data with respect to the magnetic field appears to be rather conventional and in agreement with previous reports [49]. However, a clear change appears at the structural transition temperature, below which both Hall and magnetoresistance display a much stronger field dependence. This phenomenology is compatible with a modeling within the standard Boltzmann theory of transport in which the multiband character of the electron structure, as well as the x/y anisotropy in the

nematic phase, are taken into account. Within this theoretical frame, analytical expression for both these experimental quantities can be derived. A good fit of the data is obtained in both cases using a reasonable and qualitatively consistent set of parameters. A better quantitative agreement would require the inclusion of the effect of the magnetic field on the scattering rate with the consequent need for a numerical solution of the Boltzmann equation.

The magnetoelastoresistance measured in magnetic field parallel to the crystallographic c -axis is independent of the magnetic field at any temperature. On the contrary, the magnetoelastoresistance with B_z displays a clear change at the structural transition temperature, with the absence of a magnetic-field dependence above T_S , and a more complex behavior in the nematic phase. Remarkably the two-band Boltzmann model analytically captures the independence of the magnetoelastoresistance from the magnetic field, both for in-plane fields at any temperature, and for out-of-plane fields above T_S . A good fit of the data can be obtained using a set of parameter qualitatively consistent with the ones used in the fit analysis of the magneto-resistance and Hall resistivity.

We would like to emphasize that orbital effects are usually needed to obtain a consistent picture of the transport in the nematic phase of iron-based materials, however, in our magnetoelastotransport measurements a multiband theoretical modeling appears adequate to successfully reproduce important features of the magnetic behavior of three different measured quantities. A confirmation to our hypothesis would come from the experimental analysis of other iron-based materials, exhibiting a different orbital nesting properties (e.g., BaFe_2As_2), where we expect to find similar features in the magnetoelastotransport.

DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: MULTIBAND BOLTZMANN THEORY

Here, we derive the Boltzmann equations for a multiband system in which, due to nematicity, the conductivities along the x and y directions can be different. Our starting point is the linear response equation for each component n

$$\vec{J}_n = \hat{\sigma}_n \vec{E} \quad \text{with} \quad \hat{\sigma}_n = \begin{pmatrix} \sigma_n^{xx} & \sigma_n^{xy} \\ \sigma_n^{yx} & \sigma_n^{yy} \end{pmatrix}, \quad (\text{A1})$$

where $\vec{J}_n$ is the current density of the n th band, $\vec{E}$ is the electric field, and $\hat{\sigma}_n$ is the conductivity tensor of the n th band with

components

$$\begin{aligned} \sigma_n^{\alpha\beta} &= \sigma_n^{\alpha\beta,(0)} + \sigma_n^{\alpha\beta,(1)}, \\ \sigma_n^{\alpha\beta,(0)} &= \frac{-e^2}{4\pi^3} \int d^2k v_{k,n}^\alpha v_{k,n}^\beta \tau_{k,n} \frac{\partial f_k^0}{\partial \varepsilon_{k,n}}, \\ \sigma_n^{\alpha\beta,(1)} &= \frac{e^3}{4\pi^3 \hbar} \int d^2k \frac{\partial f_k^0}{\partial \varepsilon_{k,n}} v_{k,n}^\alpha \tau_{k,n} (\vec{v}_{k,n} \times \vec{B}) \cdot \vec{\nabla}_k (v_{k,n}^\beta \tau_{k,n}), \end{aligned}$$

where $\alpha, \beta = x, y$, e is the electron charge, $v_{k,n}^\alpha$ is the α component of the group velocity of a Bloch electron with wave vector k in the n th band, $\tau_{k,n}$ is the scattering time for a Bloch electron with wave vector k in the n th band, and $f_k^0 = [e^{(\varepsilon_{k,n} - \mu)/T} + 1]^{-1}$ is the Fermi distribution function for a Bloch electron with energy $\varepsilon_{k,n}$ at a temperature T , with μ being the chemical potential. Notice that also for a nematic system $\sigma_n^{\alpha\beta,(0)} = 0$, so that only the diagonal elements of $\sigma_n^{\alpha\beta,(0)}$ contribute. Moreover, for a two-dimensional system in the xy -plane only a magnetic field along the z -direction contributes to $\sigma_n^{\alpha\beta,(1)}$, which then only contains nondiagonal elements. Finally, the Onsager reciprocity requires $\hat{\sigma}_n(B_z) = \hat{\sigma}_n^T(-B_z)$, where T implies taking the transpose matrix.

The inverse of Eq. (A1) is $\vec{E} = \hat{\rho}_n \vec{J}_n$ with

$$\begin{aligned} \hat{\rho}_n &= \hat{\sigma}_n^{-1} = \frac{1}{\sigma_n^{xx} \sigma_n^{yy} + (\sigma_n^{xy})^2} \begin{pmatrix} \sigma_n^{yy} & -\sigma_n^{xy} \\ \sigma_n^{xy} & \sigma_n^{xx} \end{pmatrix} \\ &\equiv \begin{pmatrix} \rho_n^x & B_z R_n^H \\ -B_z R_n^H & \rho_n^y \end{pmatrix}. \end{aligned} \quad (\text{A2})$$

For an N -band system we can now generalize this expression to $\vec{E} = \sum_{n=1}^N \hat{\rho}_n \vec{J}_n \equiv \hat{\rho} \vec{J}$ where $\vec{J} = \sum_n \vec{J}_n$. The expression for the full resistivity tensor is given by

$$\begin{aligned} \hat{\rho}^{-1} &= \sum_{n=1}^N \hat{\rho}_n^{-1} = \sum_{n=1}^N \hat{\sigma}_n \\ &= \sum_{n=1}^N \frac{1}{\rho_n^x \rho_n^y + (B_z R_n^H)^2} \begin{pmatrix} \rho_n^y & -B_z R_n^H \\ B_z R_n^H & \rho_n^x \end{pmatrix}, \end{aligned}$$

which, upon defining $\sigma_n^{x,y} \equiv 1/\rho_n^{x,y}$, can be written as

$$\begin{aligned} \hat{\rho}^{-1} &= \sum_{n=1}^N \frac{1}{1 + \sigma_n^x \sigma_n^y (B_z R_n^H)^2} \begin{pmatrix} \sigma_n^x & -\sigma_n^x \sigma_n^y B_z R_n^H \\ \sigma_n^x \sigma_n^y B_z R_n^H & \sigma_n^y \end{pmatrix} \\ &\equiv \begin{pmatrix} \rho^x & B_z R^H \\ -B_z R^H & \rho^y \end{pmatrix}^{-1}, \end{aligned} \quad (\text{A3})$$

that is the general definition of the resistivity in a multiband system used and discussed in the main text. For $N = 2$ bands the explicit expressions read

$$\rho^x = \frac{\sigma_1^y + \sigma_2^y + \sigma_1^y \sigma_2^y [\sigma_1^x (R_1^H)^2 + \sigma_2^x (R_2^H)^2] B_z^2}{\sigma_1^x \sigma_1^y + \sigma_2^x \sigma_2^y + \sigma_1^x \sigma_2^y + \sigma_2^x \sigma_1^y + \sigma_1^x \sigma_2^x \sigma_1^y \sigma_2^y (R_1^H + R_2^H)^2 B_z^2}, \quad (\text{A4})$$

$$\rho^y = \frac{\sigma_1^x + \sigma_2^x + \sigma_1^x \sigma_2^x [\sigma_1^y (R_1^H)^2 + \sigma_2^y (R_2^H)^2] B_z^2}{\sigma_1^x \sigma_1^y + \sigma_2^x \sigma_2^y + \sigma_1^x \sigma_2^y + \sigma_2^x \sigma_1^y + \sigma_1^x \sigma_2^x \sigma_1^y \sigma_2^y (R_1^H + R_2^H)^2 B_z^2}, \quad (\text{A5})$$

$$R^H = \frac{\sigma_1^x \sigma_1^y R_1^H + \sigma_2^x \sigma_2^y R_2^H + \sigma_1^x \sigma_1^y \sigma_2^x \sigma_2^y R_1^H R_2^H (R_1^H + R_2^H) B_z^2}{\sigma_1^x \sigma_1^y + \sigma_2^x \sigma_2^y + \sigma_1^x \sigma_2^y + \sigma_2^x \sigma_1^y + \sigma_1^x \sigma_2^x \sigma_1^y \sigma_2^y (R_1^H + R_2^H)^2 B_z^2}, \quad (\text{A6})$$

which, in the isotropic case ($\sigma_1^x = \sigma_1^y \equiv \sigma_1$, $\sigma_2^x = \sigma_2^y \equiv \sigma_2$), reduce to the well-known results

$$\rho^x = \rho^y = \frac{\sigma_1 + \sigma_2 + \sigma_1 \sigma_2 [\sigma_1 (R_1^H)^2 + \sigma_2 (R_2^H)^2] B_z^2}{(\sigma_1 + \sigma_2)^2 + \sigma_1^2 \sigma_2^2 (R_1^H + R_2^H)^2 B_z^2}, \quad (\text{A7})$$

$$R^H = \frac{\sigma_1^2 R_1^H + \sigma_2^2 R_2^H + \sigma_1^2 \sigma_2^2 R_1^H R_2^H (R_1^H + R_2^H) B_z^2}{(\sigma_1 + \sigma_2)^2 + \sigma_1^2 \sigma_2^2 (R_1^H + R_2^H)^2 B_z^2}. \quad (\text{A8})$$

In terms of the individual resistivities $\rho_{1,2}^{x,y}$, Eqs. (A4), (A5), and (A6) can be rewritten as

$$\rho^x = \frac{\rho_1^x \rho_2^x + (\rho_1^y + \rho_2^y)^{-1} [\rho_2^x (R_1^H)^2 + \rho_1^x (R_2^H)^2] B_z^2}{\rho_1^x + \rho_2^x + (\rho_1^y + \rho_2^y)^{-1} (R_1^H + R_2^H)^2 B_z^2}, \quad (\text{A9})$$

$$\rho^y = \frac{\rho_1^y \rho_2^y + (\rho_1^x + \rho_2^x)^{-1} [\rho_2^y (R_1^H)^2 + \rho_1^y (R_2^H)^2] B_z^2}{\rho_1^y + \rho_2^y + (\rho_1^x + \rho_2^x)^{-1} (R_1^H + R_2^H)^2 B_z^2}, \quad (\text{A10})$$

$$R^H = \frac{\rho_2^x \rho_2^y R_1^H + \rho_1^x \rho_1^y R_2^H + R_1^H R_2^H (R_1^H + R_2^H) B_z^2}{\rho_1^x \rho_1^y + \rho_2^x \rho_2^y + \rho_1^x \rho_2^y + \rho_2^x \rho_1^y + (R_1^H + R_2^H)^2 B_z^2}. \quad (\text{A11})$$

1. Hall coefficient

In both the isotropic and anisotropic case, the Hall coefficient can be written in the form

$$R^H = \frac{a_0 + a_1 B_z^2}{1 + a_2 B_z^2}, \quad (\text{A12})$$

where, in the isotropic case the coefficients are given by

$$a_0 = \frac{\sigma_1^2 R_1^H + \sigma_2^2 R_2^H}{(\sigma_1 + \sigma_2)^2}, \quad (\text{A13})$$

$$a_1 = \left(\frac{\sigma_1 \sigma_2}{\sigma_1 + \sigma_2} \right)^2 R_1^H R_2^H (R_1^H + R_2^H), \quad (\text{A14})$$

$$a_2 = \left(\frac{\sigma_1 \sigma_2}{\sigma_1 + \sigma_2} \right)^2 (R_1^H + R_2^H)^2, \quad (\text{A15})$$

whereas in the anisotropic case they read

$$a_0 = \frac{\sigma_1^x \sigma_1^y R_1^H + \sigma_2^x \sigma_2^y R_2^H}{\sigma_1^x \sigma_1^y + \sigma_2^x \sigma_2^y + \sigma_1^x \sigma_2^y + \sigma_2^x \sigma_1^y}, \quad (\text{A16})$$

$$a_1 = \frac{\sigma_1^x \sigma_1^y \sigma_2^x \sigma_2^y R_1^H R_2^H (R_1^H + R_2^H)}{\sigma_1^x \sigma_1^y + \sigma_2^x \sigma_2^y + \sigma_1^x \sigma_2^y + \sigma_2^x \sigma_1^y}, \quad (\text{A17})$$

$$a_2 = \frac{\sigma_1^x \sigma_2^x \sigma_1^y \sigma_2^y (R_1^H + R_2^H)^2}{\sigma_1^x \sigma_1^y + \sigma_2^x \sigma_2^y + \sigma_1^x \sigma_2^y + \sigma_2^x \sigma_1^y}. \quad (\text{A18})$$

2. Magnetoresistance

Similarly, in both the isotropic and anisotropic case, the MR can be written as

$$\frac{\Delta \rho^x(B)}{\rho^x(0)} = \frac{\rho^x(B)}{\rho^x(0)} - 1 = \frac{1 + b_0 B_z^2}{1 + b_1 B_z^2} - 1. \quad (\text{A19})$$

In the isotropic case the coefficients read

$$b_0 = \frac{\sigma_1 \sigma_2}{\sigma_1 + \sigma_2} [\sigma_1 (R_1^H)^2 + \sigma_2 (R_2^H)^2], \quad (\text{A20})$$

$$b_1 = \left(\frac{\sigma_1 \sigma_2}{\sigma_1 + \sigma_2} \right)^2 (R_1^H + R_2^H)^2, \quad (\text{A21})$$

whereas in the anisotropic case they are given by

$$b_0 = \frac{\sigma_1^y \sigma_2^y [\sigma_1^x (R_1^H)^2 + \sigma_2^x (R_2^H)^2]}{\sigma_1^y + \sigma_2^y}, \quad (\text{A22})$$

$$b_1 = \frac{\sigma_1^x \sigma_2^x \sigma_1^y \sigma_2^y (R_1^H + R_2^H)^2}{\sigma_1^x \sigma_1^y + \sigma_2^x \sigma_2^y + \sigma_1^x \sigma_2^y + \sigma_2^x \sigma_1^y}. \quad (\text{A23})$$

Note that in the latter case the coefficient b_0 is derived from $\rho^x(B_z)$, i.e., assuming that the electric field is applied along the x -direction. When the electric field is applied along the y -direction, the replacement $x(y) \rightarrow y(x)$ must be implemented in Eq. (A22).

3. Magnetoelastoresistance

We define the total resistivity anisotropy as $N = 2(\rho^x - \rho^y)/(\rho^x + \rho^y)$ and similarly for the individual components $N_n = 2(\rho_n^x - \rho_n^y)/(\rho_n^x + \rho_n^y)$. Then, from Eqs. (A4) and (A5), one obtains

$$N = 2 \frac{[\rho_2^x \rho_2^y + (R_2^H)^2 B_z^2](\rho_1^x - \rho_1^y) + [\rho_1^x \rho_1^y + (R_1^H)^2 B_z^2](\rho_2^x - \rho_2^y)}{[\rho_2^x \rho_2^y + (R_2^H)^2 B_z^2](\rho_1^x + \rho_1^y) + [\rho_1^x \rho_1^y + (R_1^H)^2 B_z^2](\rho_2^x + \rho_2^y)}. \quad (\text{A24})$$

Assuming that, in the simplest case

$$\rho_n^{x/y} = \rho_n^{x/y,0} \pm \frac{1}{2} \gamma_n \epsilon, \quad (\text{A25})$$

and defining

$$\rho_n^0 = \frac{1}{2} [\rho_n^{x,0} + \rho_n^{y,0}], \quad N_n^0 = \frac{\rho_n^{x,0} - \rho_n^{y,0}}{\rho_n^0}, \quad M_n^0 = \frac{\rho_n^{x,0} \rho_n^{y,0} + (R_n^H)^2 B^2}{(\rho_n^0)^2}, \quad (\text{A26})$$

one then finds

$$\begin{aligned} \frac{\partial N}{\partial \epsilon} = & \frac{1}{\rho_1^0 M_1^0 + \rho_2^0 M_2^0} \left(\gamma_1 \frac{\rho_2^0}{\rho_1^0} M_2^0 + \gamma_2 \frac{\rho_1^0}{\rho_2^0} M_1^0 \right) \\ & + \frac{1}{2} \frac{N_1^0 - N_2^0}{(\rho_1^0 M_1^0 + \rho_2^0 M_2^0)^2} (\gamma_1 N_1^0 \rho_2^0 M_2^0 - \gamma_2 N_2^0 \rho_1^0 M_1^0). \end{aligned} \quad (\text{A27})$$

Notice that the whole dependence of the MER on the magnetic field is contained in M_n^0 , where B appears multiplied by the n th band Hall coefficient R_n . As a consequence, one can immediately see that for in-plane magnetic field $R_n^H = 0$, so that the MER does not depend anymore on B , in agreement with our experimental findings.

In the isotropic case, using $N_n^0 = 0$, which implies $M_n = 1 + (R_n^H)^2 B^2 / (\rho_n^0)^2$, Eq. (A27) reduces to

$$\begin{aligned} \frac{\partial N}{\partial \epsilon} = & \frac{\gamma_1 \frac{\rho_2^0}{\rho_1^0} M_2^0 + \gamma_2 \frac{\rho_1^0}{\rho_2^0} M_1^0}{\rho_1^0 M_1^0 + \rho_2^0 M_2^0} \\ = & \frac{\gamma_1 \frac{\rho_2^0}{\rho_1^0} + \gamma_2 \frac{\rho_1^0}{\rho_2^0} + \frac{\gamma_1 (R_2^H)^2 + \gamma_2 (R_1^H)^2}{\rho_1^0 \rho_2^0} B_z^2}{\rho_1^0 + \rho_2^0 + \left(\frac{(R_1^H)^2}{\rho_1^0} + \frac{(R_2^H)^2}{\rho_2^0} \right) B_z^2} \approx \frac{\gamma}{\rho}, \end{aligned} \quad (\text{A28})$$

where the last step involves the approximation $\rho_n^0 \equiv \rho$, $\gamma_n \equiv \gamma$, and $R_n^H \equiv R_H$. In this limit the MER is completely independent of magnetic field. Notice, however, that since, above T_S the Hall coefficient R_H is of order $0.1 \mu\Omega/\text{T}$ and the resistivity ρ is of order $100 \mu\Omega$, the MER magnetic field dependence would be quite weak even without the above approximations, in agreement with the experimental result. Finally, we note that in the present model the temperature dependence of the coefficient $\gamma \sim 1/(T - T_S)$ has to be assumed phenomenologically, but could be, in principle, derived via a coupled Ginzburg-Landau theory.

We conclude deriving expressions suitable for the fit procedure. Let us define

$$\begin{aligned} \rho_1^0 M_1^0 + \rho_2^0 M_2^0 & \equiv K_1 + K_2 B_z^2, \\ \gamma_1 \frac{\rho_2^0}{\rho_1^0} M_2^0 + \gamma_2 \frac{\rho_1^0}{\rho_2^0} M_1^0 & \equiv K_3 + K_4 B_z^2, \\ \gamma_1 N_1^0 \rho_2^0 M_2^0 - \gamma_2 N_2^0 \rho_1^0 M_1^0 & \equiv K_5 + K_6 B_z^2, \end{aligned}$$

where

$$\begin{aligned} K_1 &= \rho_1^0 \left[1 - \frac{1}{4} (N_1^0)^2 \right] + \rho_2^0 \left[1 - \frac{1}{4} (N_2^0)^2 \right], \\ K_2 &= \left(\frac{(R_1^H)^2}{\rho_1^0} + \frac{(R_2^H)^2}{\rho_2^0} \right), \\ K_3 &= \gamma_1 \frac{\rho_2^0}{\rho_1^0} \left[1 - \frac{1}{4} (N_2^0)^2 \right] + \gamma_2 \frac{\rho_1^0}{\rho_2^0} \left[1 - \frac{1}{4} (N_1^0)^2 \right], \\ K_4 &= \frac{\gamma_1 (R_2^H)^2 + \gamma_2 (R_1^H)^2}{\rho_1^0 \rho_2^0}, \end{aligned}$$

$$\begin{aligned} K_5 &= \gamma_1 N_1^0 \rho_2^0 \left[1 - \frac{1}{4} (N_2^0)^2 \right] - \gamma_2 N_2^0 \rho_1^0 \left[1 - \frac{1}{4} (N_1^0)^2 \right], \\ K_6 &= \left(\frac{\gamma_1 N_1^0 (R_2^H)^2}{\rho_2^0} - \frac{\gamma_2 N_2^0 (R_1^H)^2}{\rho_1^0} \right), \end{aligned}$$

so that Eq. (A27) reads

$$\frac{\partial N}{\partial \epsilon} = \frac{K_3 + K_4 B_z^2}{K_1 + K_2 B_z^2} + \frac{N_1^0 - N_2^0}{2} \frac{K_5 + K_6 B_z^2}{(K_1 + K_2 B_z^2)^2}. \quad (\text{A29})$$

The fitting formula then reads

$$\frac{\partial N}{\partial \epsilon} = c_0 \frac{1 + c_1 B_z^2}{1 + c_2 B_z^2} + c_3 \frac{1 + c_4 B_z^2}{(1 + c_2 B_z^2)^2}, \quad (\text{A30})$$

with

$$\begin{aligned} c_0 &= \frac{K_3}{K_1} = \frac{\gamma_1 \frac{\rho_2^0}{\rho_1^0} \left[1 - \frac{1}{4} (N_2^0)^2 \right] + \gamma_2 \frac{\rho_1^0}{\rho_2^0} \left[1 - \frac{1}{4} (N_1^0)^2 \right]}{\rho_1^0 \left[1 - \frac{1}{4} (N_1^0)^2 \right] + \rho_2^0 \left[1 - \frac{1}{4} (N_2^0)^2 \right]}, \\ c_1 &= \frac{K_4}{K_3} = \frac{\frac{\gamma_1 (R_2^H)^2 + \gamma_2 (R_1^H)^2}{\rho_1^0 \rho_2^0}}{\gamma_1 \frac{\rho_2^0}{\rho_1^0} \left[1 - \frac{1}{4} (N_2^0)^2 \right] + \gamma_2 \frac{\rho_1^0}{\rho_2^0} \left[1 - \frac{1}{4} (N_1^0)^2 \right]}, \\ c_2 &= \frac{K_2}{K_1} = \frac{\frac{(R_1^H)^2}{\rho_1^0} + \frac{(R_2^H)^2}{\rho_2^0}}{\rho_1^0 \left[1 - \frac{1}{4} (N_1^0)^2 \right] + \rho_2^0 \left[1 - \frac{1}{4} (N_2^0)^2 \right]}, \\ c_3 &= \frac{N_1^0 - N_2^0}{2} \frac{K_5}{K_1} \\ &= \frac{N_1^0 - N_2^0}{2} \frac{\gamma_1 N_1^0 \rho_2^0 \left[1 - \frac{1}{4} (N_2^0)^2 \right] - \gamma_2 N_2^0 \rho_1^0 \left[1 - \frac{1}{4} (N_1^0)^2 \right]}{\left\{ \rho_1^0 \left[1 - \frac{1}{4} (N_1^0)^2 \right] + \rho_2^0 \left[1 - \frac{1}{4} (N_2^0)^2 \right] \right\}^2}, \\ c_4 &= \frac{K_6}{K_5} = \frac{\left(\gamma_1 N_1^0 \frac{(R_2^H)^2}{\rho_2^0} - \gamma_2 N_2^0 \frac{(R_1^H)^2}{\rho_1^0} \right)}{\gamma_1 N_1^0 \rho_2^0 \left[1 - \frac{1}{4} (N_2^0)^2 \right] - \gamma_2 N_2^0 \rho_1^0 \left[1 - \frac{1}{4} (N_1^0)^2 \right]}. \end{aligned}$$

Notice that if we set $N_n^0 = 0$, the expression reduces to

$$\frac{\partial N}{\partial \epsilon} = c_0 \frac{1 + c_1 B_z^2}{1 + c_2 B_z^2}, \quad (\text{A31})$$

with fitting parameters given by

$$c_0 = \frac{1}{\rho_1^0 + \rho_2^0} \frac{1}{\rho_1^0 \rho_2^0} [\gamma_1 (\rho_2^0)^2 + \gamma_2 (\rho_1^0)^2], \quad (\text{A32})$$

$$c_1 = \frac{\gamma_1 (R_2^H)^2 + \gamma_2 (R_1^H)^2}{\gamma_1 (\rho_2^0)^2 + \gamma_2 (\rho_1^0)^2}, \quad (\text{A33})$$

$$c_2 = \frac{1}{\rho_1^0 + \rho_2^0} \left(\frac{(R_1^H)^2}{\rho_1^0} + \frac{(R_2^H)^2}{\rho_2^0} \right). \quad (\text{A34})$$

The assumption $N_n^0 = 0$ means that $\rho_n^{x,0} - \rho_n^{y,0} = 0$. This is certainly true in the isotropic case, however, it can also be assumed in the nematic phase if we imply that the anisotropy is proportional to the strain only, see Eq. (A25).

TABLE I. Parameters for fitting the Hall coefficient Eq. (A12) to experimental data.

T [K]	a_0 [$\mu\Omega$ cm/T]	a_1 [$\mu\Omega$ cm/T ³]	a_2 [1/T ²]
15	-0.3931	0.0029	0.02571
20	-0.466	0.00135	0.0271
30	-0.3191	0.00159	0.01638
40	-0.1239	0.00163	0.00976
50	-0.057	0.00086	0.00868
60	0.02423	0.000456	0.003967
80	0.0823	-1.265×10^{-5}	7.698×10^{-5}
90	0.0556	-1.51×10^{-5}	1.874×10^{-4}
120	0.0269	-0.00216	0.30396

APPENDIX B: FIT OF THE EXPERIMENTAL RESULTS

In this Appendix we summarize the results of the fit procedure for Hall resistivity, MR and MER, using Eqs. (A12), (A19), and (A30).

In the main text, we show in Fig. 2(b) the fit of the experimental data for the Hall resistivity using Eq. (A12). As one can see, the procedure yields a rather good fit to the data for $\rho_{xy} = R^H B_z$. The fitting parameters are given in Table I.

The analysis of the MR is shown in Fig. 2(a), where the theoretical fit is shown on top of the experimental data using Eq. (A19). The corresponding fitting parameters are listed in Table II. As we already remarked in the main text, within the present derivation, the parameters a_2 and b_1 describe the same parameter combination. However, the fit procedure lead to a difference of approximately a factor of 7.0–8.0 at low temperature. In principle, one can simultaneously fit the Hall resistivity and the MR, constraining $b_1 = a_2$, finding the results shown in Table III. The overall trend of the fitting parameters is maintained, although the introduction of a constraint worsens the fit, especially at low temperature, as it can be seen in Fig. 4. However, it must be pointed out that the parameters b_1 and a_2 are only relevant for fitting the MR and ρ^{xy} at large fields. The discrepancy in the unconstrained fit may then point toward a violation of the standard perturbative treatment of the Boltzmann equation, as discussed in the main text.

The analysis of the MER is shown in Fig. 3(c), where we show the experimental data together with the theoretical polynomial fits Eq. (A30). From these fits we extract the magnetic field dependence at different temperatures. The corresponding curves are presented in Fig. 5, which shows a positive trend

TABLE II. Parameters for fitting the MR Eq. (A19) to experimental data.

T [K]	b_0 [1/T ²]	b_1 [1/T ²]
15	0.01078	0.00364
20	0.00922	0.003495
30	0.00673	0.00333
40	0.00636	0.00407
50	0.00346	0.00231
60	0.01195	0.01063
90	0.28413	0.27859

TABLE III. Parameters for fitting the Hall coefficient Eq. (A12) the MR Eq. (A19) to experimental data, with the constraint $b_1 = a_2$.

T [K]	a_0 [$\mu\Omega$ cm/T]	a_1 [$\mu\Omega$ cm/T ³]	b_0 [1/T ²]	$b_1 = a_2$ [1/T ²]
15	-0.277695	0.00197007	0.0215466	0.01
20	-0.340067	0.00159451	0.0161647	0.008
30	-0.262895	0.00139701	0.013011	0.008
40	-0.110791	0.00140892	0.00992253	0.007
50	-0.0463063	0.00066924	0.00650638	0.005
60	0.020013	0.00066443	0.00804311	0.007
90	0.0653995	0.00769684	0.153141	0.15

for all temperatures and an enhanced field dependence in the temperature range above T_C where superconducting fluctuations are expected to be relevant.

Upon fitting the MER using Eq. (A30), one realizes that the inclusion of coefficients c_3 and c_4 does not improve the fit. Alternatively, one could also do the fitting with only coefficients c_2 , c_3 , c_4 with a similar accuracy, however, the first choice implies that $N_n^0 \approx 0$ whereas the second choice is harder to justify. Let us then assume $N_n^0 \approx 0$ so that the MER can be written as Eq. (A31). In this approximation, we find the set shown in Table IV. We use these coefficients to plot the temperature-dependent curves shown in Fig. 3(c). The agreement between the theoretical curves and the data is not to be interpreted as an independent fit of the temperature dependence, rather as a representation of the consistency of the field-dependent analysis across the measured temperature range.

Figure 6 shows the coefficients extracted from the Hall analysis, Eq. (A12). Figure 6(a) reports a_0 , associated with the low-field Hall response. In Fig. 6(b), instead of showing a_1 and a_2 separately, we plot the combination a_2/a_1 , which in the two-band model corresponds to the quantity $1/R_1^H + 1/R_2^H$. For comparison, Fig. 6(b) also shows the ratio b_1/a_1 extracted from the MR fit, Eq. (A19), under the constraint $b_1 = a_2$. The agreement is good over the temperature range where both Hall and MR signals are sizable. At higher temperatures,

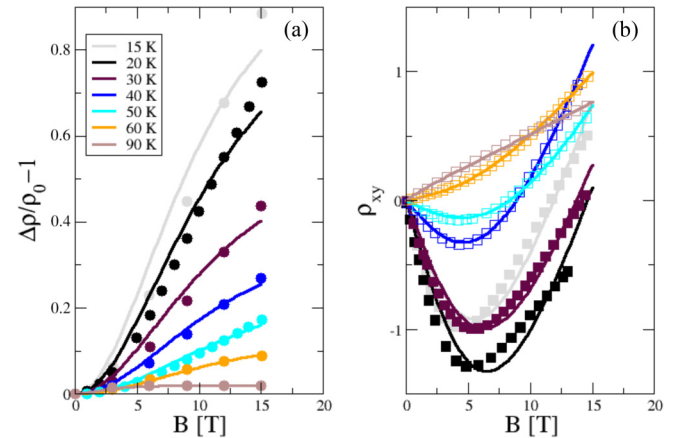


FIG. 4. Simultaneous fit of the MR and of the Hall resistivity, with the constraint $b_1 = a_2$. In both panels (a) and (b), the theoretical fit using Eqs. (3) and (2) are shown as solid lines on top of the experimental data (symbols).

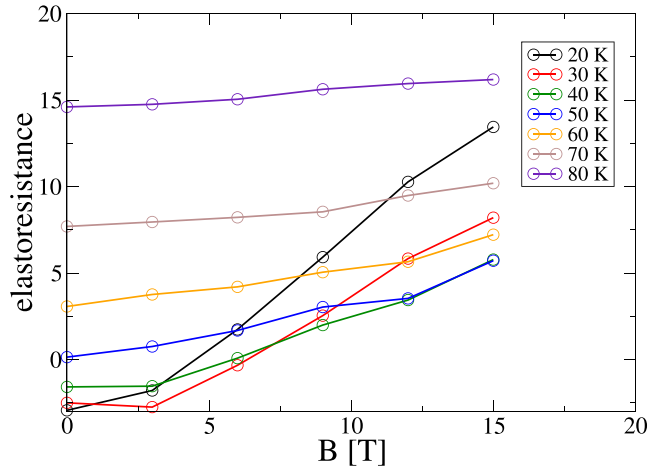


FIG. 5. Magnetic-field dependence of the MER extracted from the polynomial fits at various temperatures below T_S .

however, the fit parameters must be interpreted with caution. Above 60 K the MR becomes very small, which according to Eq. (A19) can be obtained for any $b_0 = b_1$, while ρ_{xy} becomes nearly linear, implying $a_1 \sim a_2 \sim 0$, and eventually very small, so that also $a_0 \sim 0$. As a consequence, the individual fitting parameters are only weakly constrained in this regime.

Figure 7 shows the temperature dependence of the coefficients extracted from the MER fit, Eq. (A30). Notice that the sharp variation of c_1 we observe here is not anomalous, but follows directly from its definition: according to Eqs. (A32) and (A33), c_1 develops a pole when c_0 crosses zero. We also note if γ_i and ρ_i for the two bands are the same then $c_1 = c_2$ which seems to be the case for temperatures close to T_S .

TABLE IV. Parameters for fitting the MER Eq. (A31) to experimental data.

T [K]	c_0	c_1 [1/T ²]	c_2 [1/T ²]
20	-3.037	-0.0467	0.005
25	-3.2	-0.032	0.004
30	-2.991	-0.0255	0.003
35	-2.476	-0.024	0.0028
40	-1.724	-0.0028	0.0026
45	-0.775	-0.052	0.0025
50	0.367	0.096	0.0023
55	1.73	0.0177	0.0018
60	3.36	0.0073	0.0011
65	5.33	0.0034	0.0005
70	7.76	0.0018	0.0003
75	10.8	0.002	0.0011
80	14.5821	0.0087	0.00742

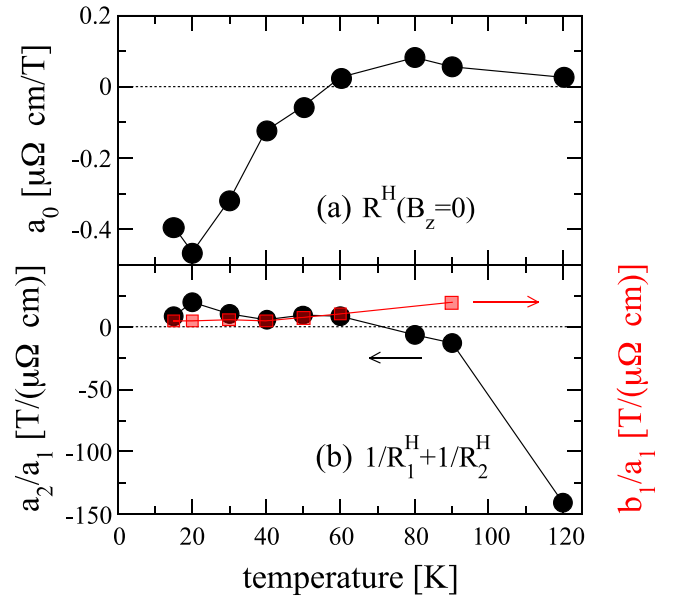


FIG. 6. Temperature dependence of the coefficients extracted from the Hall fit Eq. (A12). Panel (a) shows the coefficient a_0 . Panel (b) shows the combination a_2/a_1 , together with the corresponding quantity extracted from the MR fit Eq. (A19). The comparison is meaningful in the temperature range where both Hall and MR signals are sizable, while at higher temperature the vanishing MR and the nearly linear, very small Hall response make the corresponding fit parameters weakly constrained.

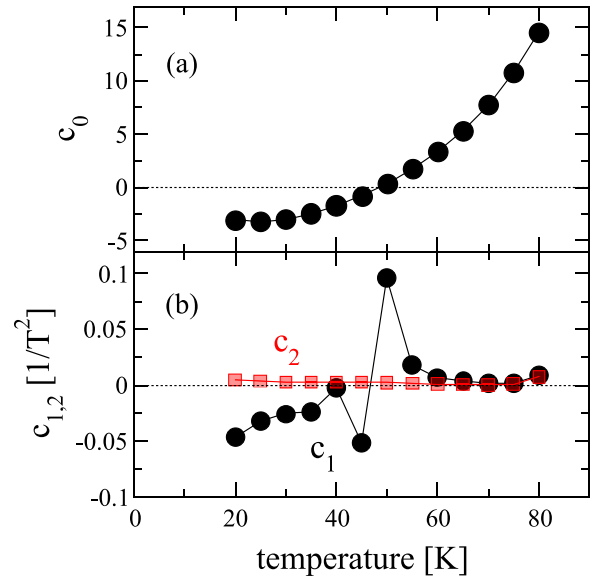


FIG. 7. Temperature dependence of the coefficients extracted from the MER fit Eq. (A30). Panel (a) shows c_0 , while panel (b) reports c_1 and c_2 .

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