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Effect of the screened Coulomb disorder on magneto-transport in Weyl semimetals

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The observation of negative longitudinal magneto-resistivity (NLMR) in Weyl semimetals has gained strong support in recent experiments. It is believed that charged impurities play an important role in the measurement of NLMR. We thus employ a screened Coulomb disorder to model charged impurities and derive a general screening length depending on the magnetic field, chemical potential and temperature. We study the magneto-transport in a two-node Weyl semimetal in which the intra-valley scattering and the inter-valley scattering can be explored simultaneously. We also calculate the effect of the misalignment of the external electric field and the magnetic field on the longitudinal and transverse magnetoconductivities, recovering the experimental observations. We show that the former (latter) is suppressed (enhanced) sensitively with the density of the impurity. This feature makes it hard to observe the NLMR in experiments in the heavy doping case. These results may be exploited to explain the sample-dependent observation of NLMR and deepen our understanding of magneto-transport in Weyl semimetals. Published by AIP Publishing.

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I. INTRODUCTION

Weyl semimetal, a new kind of topological material, proposed a few years ago,1 has drawn much attention. Two kinds of Weyl semimetals have been identified in the TaAs family,2–5 YbMnBi2,6 and Mo,W1–xTe2.7–9 The band crossings named Weyl nodes behave like monopoles due to the Berry curvature in the three-dimensional (3D) momentum space. The Weyl nodes appear in pairs with opposite chirality, equipped with a property called chiral anomaly.10 Anomalous transport caused by the chiral anomaly is a vital process in condensed matter physics as well as the cases in high energy physics and astrophysics. This may manifest itself as negative longitudinal magneto-resistivity (NLMR) (or positive magneto-conductance). The positive longitudinal magneto-conductance has recently been observed in Na3Bi,11 TaAs,12,13 NbP,14 GdPtBi,15 Cd3As2,16 and TaP.17 The mixed axial gravitational anomaly in NbP has also been observed recently,18 which may open a new window in theoretical aspect.

Although there have been intensive investigations,19–24 they are still insufficient to understand the realistic features observed in experiments because most of the previous studies dealt with a one-node model. It is known that the Weyl nodes exist in pairs and the chiral anomaly involves at least one pair. Therefore, it is not possible to explore the chiral anomaly in such a one-node model.10,25 Furthermore, it is noted that the observed magneto-resistance in some experiments is not negative but positive;26 while it is negative in some others. Another observation is that the Fermi level varies dramatically in different samples. Changing the location of the Fermi level in experiments can be realized by introducing donors and acceptors at different concentrations. These sensitively sample-dependent features indicate that the charged impurities may be the main type of disorder in transport measurements. Here, we employ a Coulomb screening disorder to describe the effect of the charged impurities, and the strength of the screening effect is encoded in the screening length. In a previous work,27 to study the contribution of inter-valley scattering to the magnetoconductivity explicitly, we fixed the screening length in each numerical evaluation. In a more general case, the screening length may depend on the magnetic field strength, the chemical potential and the temperature. We therefore study this general situation and give a deeper insight on the exploration of chiral magneto-transport.

We showed that NLMR can appear in a magnetic field at low temperatures and in the light doping case, consistent with the prediction of chiral anomaly.22 We also studied the effect of temperature and the density of impurity. We found that the NLMR vanishes thoroughly with the increase of impurity, which is attributed to the cover up effect of transverse magnetoconductivity. This effect is mainly caused by
different dependences of the transverse and longitudinal magnetoconductivities on the impurity scattering matrix elements. In realistic measurements, it is not easy to parallelize the magnetic field and the electric field exactly, and then the transverse resistivity may be mixed with a longitudinal resistivity. We checked this effect through the angular dependence of the magneto-conductivity. Based on these results, we can understand recent observations in experiments clearly and propose a way to facilitate the observation of the NLMR in future experiments. We also showed that the inter-valley scattering can be as important as the intra-valley scattering, in future experiments. We also showed that the inter-valley scattering can be as important as the intra-valley scattering, even more vital than the latter. The results of inter-valley scattering are novel since previously the one-node model was not considered for the inter-valley scattering.

This paper is organized as follows. In Sec. II, we study a two-node model for a Weyl semimetal and show the expression for the Coulomb screening length and the formulas for the longitudinal and transverse magnetoconductivities in the presence of the Coulomb screening disorder are also given. In Sec. III, we study the effect of the magnetic field, the temperature and the density of impurity on the longitudinal and transverse magnetoconductivities, respectively. The angular dependence of magneto-conductivity is studied in Sec. IV. We further interpret the vanishing of the NLMR in Sec. V. The conclusions are given in Sec. VI.

II. TWO-NODE MODEL AND FORMALISM

In a two-node model for a Weyl semimetal, the energy dispersion is linear in momentum around the nodes forming cones (or valley), labeled \( K \) and \( K' \). In the absence of an external magnetic field \( \mathbf{B} \), the Hamiltonian for the two-node model \(^{28} \) is given by

\[
\mathcal{H} = \hbar v_F \left[ (\mathbf{k} + \nu \mathbf{k}_c) \cdot \sigma \right],
\]

where \( \nu = 1(-1) \) for the \( K (K') \) cone, \( v_F \) is the Fermi velocity (we take this value as \( v_F = 10^6 \text{ m/s} \)) and \( \mathbf{k}_c = (0, 0, \pm k_z) \) represents the positions of the two Weyl nodes in the momentum space.

When a magnetic field \( \mathbf{B} = (0, 0, B) \) is applied along the \( z \) direction, we have

\[
\mathbf{k} = \left( k_x - \frac{eB}{\hbar} y, -i\partial_y, k_z \right),
\]

under the Landau gauge \( \mathbf{A} = (-yB, 0, 0) \). With the help of ladder operators

\[
\begin{align*}
\lambda &= -\left[(y - i\partial_y)k_z / l_B + i\partial_y\right] / \sqrt{2}, \\
\lambda^\dagger &= -\left[(y - i\partial_y)k_z / l_B - i\partial_y\right] / \sqrt{2},
\end{align*}
\]

we obtain the Hamiltonian in the \( K \) cone

\[
H_K = \hbar v_F \begin{pmatrix}
-k_z & -\frac{\sqrt{2}}{l_B} \lambda \\
\frac{\sqrt{2}}{l_B} \lambda^\dagger & -k_z
\end{pmatrix} - \hbar v_F \mathbf{k}_c \cdot \sigma,
\]

where \( l_B = \sqrt{\hbar/eB} \) is the magnetic length.

As the Fermi level leaves the Weyl nodes and does not reach the first Landau band, the semiclassical conductivity of the 0-th Landau band can be derived as \( \sigma_{zz} = e^2 n_F D \), where \( n_F = (1/2\pi\alpha) (2/\pi \hbar v_F) \), \( D = v_F^0 t_{0,0} \) is the diffusion coefficient, \( t_{0,0} \) is the transport time, and \( k_F \) is the Fermi wave number. Under the first-order Born approximation, for the scattering among the states on the Fermi surface of the 0-th band, we obtain the relaxation time by

\[
\tau_{0,0} = 2\pi \sum_{k'_{ij}} \langle \langle \mathcal{U}_{k,i,k',j}^{0,0} \rangle \rangle \Delta \delta(E_F - E_{k'}^0) \times \left( 1 - v_{0,k'}^r / v_F \right),
\]

where \( E_F \) is the Fermi energy, which is derived below, and \( E_{k'}^0 \) is the eigenenergy of the 0-th Landau level. Here, \( \mathcal{U}_{k,i,k',j}^{0,0} \) represents the scattering process between the eigen states \( |0, k_x, k_z \rangle \) and the eigen state \( |0, k'_x, k'_z \rangle \) and \( \langle \langle \mathcal{U}_{k,i,k',j}^{0,0} \rangle \rangle^2 \) is the corresponding scattering matrix elements. \( v_{0,k'}^r \) is the velocity of an electron along the \( z \) direction with the wave number \( k'_z \) in the 0-th Landau band \(^{22} \) and \( \Delta \) is introduced to correct the unphysical van Hove singularity near the band edge. \(^{22} \) Other details can be found in our previous work. \(^{27} \) The screened Coulomb potential with charged impurities is

\[
\phi(r) = \frac{e^2}{\varepsilon_0 r^*} e^{-r^* / \lambda},
\]

where \( \varepsilon_0 \) represents the dielectric constant, \( \alpha \) is the dimensionless constant, \( e \) is the charge of electrons, and \( \lambda \) is the screening length.

In an undoped Weyl semimetal, the Fermi level is located at the Weyl nodes. Adding donor and acceptor impurities will introduce extra charges to the energy bands, which shift the Fermi level away from the Weyl nodes. Therefore, the excess electron density \( n_{imp}(\mu, T) \) is given by

\[
n_{imp}(\mu, T) = n_D - n_A,
\]

where \( n_D \) and \( n_A \) represent the density of the donor and the acceptor, respectively. And, the excess electron density is calculated via

\[
n_{imp}(\mu, T) = g \sum_{k_z} \left[ f_0(\nu \hbar v_F k_z - \mu) - f_0(\nu \hbar v_F k_z) \right]
\]

\[
= 2g \sqrt{\frac{1}{\pi}} \hbar v_F k_B T \ln \left( \frac{e^{\beta_0} + 1}{2} \right),
\]

where \( f_0(x) = \frac{1}{1 + e^{\beta x}} \) is the Fermi function and \( g \) represents the spin and Landau level degeneracy. The summation over \( k_z \) in Eq. (8) will be finally turned into an integration convention, which is from \( -\infty \) to \( \infty \) according to Ref. 30. More specifically, the integration in Eq. (8) can be done for two cones as four parts: (1) \( \nu = +, k_z > 0 \) represents the electron carrier in the \( K \) cone; (2) \( \nu = -, k_z < 0 \) represents the electron carrier in the \( K' \) cone; (3) \( \nu = +, k_z < 0 \) represents the hole carrier in the \( K \) cone; and (4) \( \nu = -, k_z > 0 \) indicates the hole carrier in the \( K' \) cone.
In experiments, $n_{\text{imp}}$ is fixed for every sample with donors and acceptors. Therefore, by taking $\mu = E_F$ in Eq. (8), we can obtain $E_F$

$$E_F \equiv k_B T \ln \left( 2 e^{\frac{\mu}{k_B T^*}} - 1 \right).$$  \hspace{1cm} (9)

where $\lambda$ can be determined by the Thomas-Fermi approximation

$$\frac{1}{\lambda^2} = \frac{4 e^2}{\pi e_0} \frac{\partial n(\mu, T)}{\partial \mu} = \frac{4 e^2}{\pi e_0} \times \frac{4 e^2 v_F^2}{\hbar^2}(e^2 v_F^2 + 1).$$  \hspace{1cm} (10)

Under a magnetic field, the degeneracy of each Landau level per unit area (each $k_z$) is given by $g_0 = \frac{e B}{2 \pi m^*}$, then $g = 2 \times g_0$, where 2 comes from the spin degeneracy. Substituting this into Eq. (10), we obtain

$$\frac{1}{\lambda^2} = \frac{2 e^3 B}{\pi e_0 \hbar^2 v_F} \times \frac{8 e^2 v_F^2}{e^2 v_F^2 + 1}.$$  \hspace{1cm} (11)

As the Fermi energy is located at the Weyl nodes, i.e., $\mu = 0$, we recover the formula obtained in Ref. 30 up to a constant, i.e., 4 which takes the spin and Weyl node degeneracy. Therefore, our result is more general than the result derived previously and can give correct dependence of the screening length on various external parameters.

In Fig. 1, we show the screening length and the magnetic length versus the magnetic field strength at $T = 1$ K. It is seen that the screening length is about 2.5 nm under a weak magnetic field and decreases slightly at larger $B$. However, the screening length is shorter than the magnetic length in all the range of the studied magnetic field strength.

The longitudinal conductivity $\sigma_{zz}$ can be derived by the expression

$$\sigma_{zz} = \frac{e^2 h}{2 \pi v F} \sum_{k_z, k_y} \text{Tr} \left( \mathbf{v}_{0, k_z}^\prime \mathbf{G}_{0, k_z}^R \mathbf{v}_{0, k_z}^\prime \mathbf{G}_{0, k_z}^A \right).$$  \hspace{1cm} (12)

where the bold symbols, e.g., $\mathbf{v}_{0, k_z}^\prime$, $\mathbf{v}_{0, k_z}^\prime$, and $\mathbf{G}_{0, k_z}^R$, are all 2 \times 2 diagonal matrices in the valley space (i.e., cone space), and the subscript 0 represents the $n = 0$ Landau subspace. $\mathbf{v}_{0, k_z}^\prime$ is the corresponding velocity mentioned above, $\mathbf{v}_{0, k_z}^\prime$ is the dressed velocity after taking into account the vertex correction and $\mathbf{G}_{0, k_z}^R$ is the retarded/advanced Green function, which can be written explicitly as

$$\mathbf{v}_{0, k_z}^\prime = \begin{pmatrix} v_{0, k_z}^\prime & 0 \\ 0 & v_{0, k_z}^\prime \end{pmatrix},$$

and $i$ represents $K$ (corresponding to $\nu = +$) and $K'$ (corresponding to $\nu = -$) cones.

The retarded/advanced Green’s function for the 0-th band in the $K$ cone is $\mathbf{G}_{0, k_z}^R = 1/(E_F - E_{0, k_z} + i\hbar/2\tau_{0, k_z}^\text{tr})$, where $\tau_{0, k_z}^\text{tr}$ is the corresponding transport time which is related to the transition probability (induced by the impurities) in the $K$ cone at $n = 0$ band. In the $K'$ cone, a similar expression of Green’s function can be obtained.

Substituting Eq. (13) into Eq. (12), we have

$$\sigma_{zz} = \frac{e^2 h}{2 \pi v F} \sum_{k_z, k_y} \left( \mathbf{v}_{0, k_z}^\prime \mathbf{G}_{0, k_z}^{R/A} \mathbf{G}_{0, k_z}^{R/A} \right).$$  \hspace{1cm} (15)

In the diffusive regime, one can replace $\mathbf{G}_{0, k_z}^{R/A}$ by $(2\pi/\hbar)\tau_{0, k_z}^\text{tr} \delta(E_F - E_{0, k_z}^\text{tr})$.\hspace{1cm} (16)

Expanding the $\delta$-function, we obtain

$$\sigma_{zz} = \frac{e^2}{V} \sum_{k_z, k_y} \left( \mathbf{v}_{0, k_z}^\prime G_{0, k_z}^{R/A} \right)^2 \Lambda_{k_y}^{0, k_z} \delta(E_F - E_{0, k_z}^\text{tr})$$

$$= \frac{e^2 v_F}{\hbar L_x L_y} \sum_{k_z} \tau_{k_y}^{0, k_z} \Lambda,$$

where $V = L_x L_y L_z$ is the volume of a box. We note that $\sigma_{zz}$ is proportional to the transport time, which is inversely proportional to the scattering matrix probability [Eqs. (5) and (17)].

This transport time dependence is also confirmed in Ref. 22. After a straightforward calculation which is shown in S2 in Ref. 34, we obtain
the scattering matrix probability. This is opposite to that of longitudinal magneto-conductivity, which can give rise to interesting transport phenomena to be studied below. Other expressions appearing in Eq. (20) can be found in S5 in Ref. 34.

III. LONGITUDINAL AND TRANSVERSE CONDUCTIVITIES

We show the magneto-conductivity as a function of magnetic field in Fig. 2. It is seen that the dependence of the transverse magnetoconductivity (σxx) on the magnetic field is distinct from that of the longitudinal magnetoconductivity (σzz). σxx decreases, while σzz increases with B. To understand this feature, we note that the orbits of electrons in a 3D space in the presence of a magnetic field fall on a cylindrical surface in which the projection on the x–y plane is a circle; while the axis of the cylinder is along the z-direction. Therefore, σxx involves the hopping between neighbors of these cylindrical orbits, which describes the transport of carriers along the x-direction driven by a field in the same direction. Based on the hopping feature, the transport encoded by σxx behaves like that in an insulator. For a stronger magnetic field, the radius of the cylinder is shrunk and it becomes harder to realize the hopping, leading to decreased conductivity observed in Fig. 2(a). In addition, we observe that the inter-valley contribution to the conductivity is smaller than that of intra-valley contribution when B < 5 T. For larger B, these two contributions tend to be approximate. For σzz, the transport along the z-direction occurs on the surface of cylindrical orbits. Therefore, each cylinder can be regarded as a conducting channel. In this sense, the transport along the z-direction resembles that occurs in a metal. Under a stronger magnetic field, the density of the cylinders is increased so that the number of channels is increased, giving rise to a larger conductivity. We further evaluate the density of the magnetic flux (or degeneracy per energy level) as $2\pi l_B^2$. Therefore, the conductivity should be proportional to $B$ in this picture. However, a $B^2$-dependence of the longitudinal magneto-resistivity was observed in experiments. 12,26 And, the NLMR is explained in view of the chiral
anomaly.\textsuperscript{19,20} Although the results from the conducting-channel picture are not quantitative to the observations, it might exist in principle. In a strong magnetic field, the chiral anomaly may be dominant, while the conducting-channel picture may be dominant in a weaker magnetic field. Nevertheless, the conducting-channel picture is instructive to understand the behavior of $\sigma_{zz}$ qualitatively.

To understand the behavior of magneto-transport deeply, we study the effects of the temperature and the density of the impurity, which play important roles in transport. For example, in Ref. 12, the NLMR is temperature dependent and is observed only at low temperatures. We thus show $\sigma_{xx}$ and $\sigma_{zz}$ as a function of temperature in Fig. 3. For $\sigma_{xx}$, both the intra-valley and inter-valley processes enhance the conductivity at higher temperatures, but with different slopes. $\sigma_{zz}$ decreases with increasing temperature. The distinct features stem from different dependences of conductivities on $\theta_i^2$. That means the intra-valley conductivities ($\sigma_{xx, intra}$ and $\sigma_{zz, intra}$) are determined by $\cos^2(\theta_i^2/2)$. This is just opposite to that of intra-valley conductivities ($\sigma_{xx, inter}$ and $\sigma_{zz, inter}$) which is related to $\sin^2(\theta_i^2/2)$ (see Eq. (S37) in Ref. 34). $\theta_i^2$ is proportional to $E_F$ (due to the dependence on $k_z$ or $k_x$) which increases at higher temperatures [see Eq. (9)]. Hence, $\cos^2(\theta_i^2/2)/\sin^2(\theta_i^2/2)$ decreases (increases) with temperature, so does $\sigma_{xx, inter}/\sigma_{xx, intra}$. $\sigma_{zz}$ monotonically decreases with temperature in Fig. 3(b), which may lead to the NLMR at higher temperatures.

As disorder plays a crucial role in the magneto-transport, we show its effect on $\sigma_{xx}$ and $\sigma_{zz}$ in Fig. 4, which can be understood in view of the physical picture of the transport in x and z directions described above. For the transport in the x-y plane, increasing the density of the impurity enhances the hopping, leading to a larger $\sigma_{xx}$, as observed from Fig. 4(a). This is particularly valid for $\sigma_{xx}^{K}$ since only intra-valley scattering is involved. $\sigma_{zz}$ only grows gradually with the density of impurity. This may be a consequence of inter-valley scattering with large momentum changes. Recall that the two valleys are located along the z direction. Therefore, $\sigma_{xx}^{K}$ describes the transport in the x-y plane with the aid of transition along the z direction. Thus, it only grows slowly with increasing density of impurity due to such a secondary involved process. For $\sigma_{zz}$, transport along the z direction occurs on the surface of the cylinder so that the impurity scattering would prevent the chiral current in the z-direction, giving rise to the decreasing behavior observed in Fig. 4(b). This is analogous to the behavior of impurity in a metal.

![Fig. 3](image3.png)

**FIG. 3.** (a) Same as Fig. 2 but as a function of $T$. (b) The longitudinal magnetoconductivity $\sigma_{zz}$ as a function of $T$. Here, $n_{imp} = 1 \times 10^{15}$ cm$^{-3}$ and $B = 1$ T.

![Fig. 4](image4.png)

**FIG. 4.** (a) Same as Fig. 2, but as a function of the density of impurity $n_{imp}$. (b) The longitudinal magnetoconductivity $\sigma_{zz}$ as a function of the density of impurity $n_{imp}$. Here, $T = 1$ K and $B = 1$ T.

### IV. ANGULAR DEPENDENCE OF MAGNETO-CONDUCTIVITY

In a recent experiment,\textsuperscript{18} the angular dependence of the magneto-conductivity was measured, which is defined as $\Delta \sigma_{zz} = \sigma_{zz}(B \neq 0) - \sigma_{zz}(B = 0) = \cos^2 \phi \sigma_{zz} + \sin^2 \phi \sigma_{xx} - \sigma_{zz}^{D}$, where $\phi$ is the misalignment angle of the electric field $E$ from the magnetic field $B$ and $\sigma_{zz}^{D}$ is the corresponding reference value. Under this disposal, $\sigma_{zz}$ and $\sigma_{xx}$ are combined with each other at different weights. Note from Eqs. (18) and (S33) in Ref. 34 that the dielectric constant $\varepsilon$ appears in the nominator (denominator) for $\varepsilon_{xx}(\varepsilon_{zz})$, hence this subtle orientation dependence of the dielectric constants needs to be considered in numerical calculations. In Ref. 33, the dielectric constants for TaAs are calculated as $\varepsilon_{xx} = 871$ and $\varepsilon_{zz} = 36$ in the long-wavelength limit via density functional theory (DFT). We further calculated the static dielectric constants of TaAs using DFT (see the S6 section in Ref. 34). A similar figure to that in Ref. 18 is obtained in Fig. 5(a). We note that when $\phi$ is close to $0^\circ$ (E $\parallel$ B), a positive longitudinal magneto-conductivity (corresponding to the NLMR) is recovered. This is just the case described by the chiral anomaly. When $\phi$ is close to $90^\circ$ (corresponding to the E $\perp$ B), the measured magneto-conductivity becomes negative (which means the vanishing of the NLMR). The combinations of $\sigma_{xx}$ and $\sigma_{zz}$ may suppress the observation of the NLMR of $\sigma_{zz}$. It can be concluded that the observed magneto-conductivity is sensitive to the misalignments of $E$ and $B$.

Furthermore, we show the effect of impurity on the angular dependence of the magneto-conductivity in Fig. 5(b). When the density of impurity is low (light doping), the NLMR can be observed in a wide range of $\phi$, except for the
neighbors of ±90°. When increasing the density of impurity, the NLMR is suppressed dramatically. At the same time, the range of φ in which the NLMR can be observed becomes narrow. This means that the observation of the NLMR becomes more difficult in the heavy doping case. It should be emphasized that this degradation of the NLMR due to the presence of disorder can also be derived by using an isotropic dielectric constant.

V. VANISHING NLMR AROUND E \parallel B

It is noted that the NLMR (i.e., positive magneto-conductivity \( \sigma_{zz} \)) measured at \( E \parallel B \) was not observed in every experiment.\(^{26} \) And, in some experiments,\(^{12} \) it was observed only as \( B > 1 \) T. A reasonable explanation is that the transverse resistivity (related to \( \sigma_{xx} \)) may be mixed with a longitudinal resistivity in realistic measurements since it is not easy to parallelize the magnetic field and the electric field exactly. A perpendicular component of the electric field (e.g., \( E_z \)) with respect to the magnetic field may exist, inducing a large transverse magneto-resistivity (MR), which may overwhelm a smaller longitudinal MR.

In order to reflect the realistic measurement in the experiment, we define \( \sigma_{zz}^{mix} = \eta \sigma_{zz} + (1 - \eta) \sigma_{xx} \) to describe the case. We take \( \eta = 0.95 \) (1 - \( \eta = 0.05 \) characterizes the misalignment of \( E \) from \( B \)) in our calculation. We plot the mixed longitudinal magneto-conductivity at different densities of impurity. In Fig. 6(a) (corresponding to light doping), positive \( \sigma_{zz}^{mix} \) (i.e., the NLMR) still appears; while in Fig. 6(b) (corresponding to heavy doping), \( \sigma_{zz}^{mix} \) decreases with increasing \( B \) to a U-turn point at about 0.75 T. This is just a negative longitudinal magneto-conductivity, which means a positive longitudinal magneto-resistivity. After this U-turn point, the longitudinal magneto-conductivity rises with \( B \). Therefore, the observation of the NLMR is sensitive to disorder, and is confined to limited cases.

VI. SUMMARY

In this work, we studied the longitudinal and transverse magneto-conductivities of Weyl semimetals in the presence of screened Coulomb disorder via a two-node model. We derived a general formula for the screening length which depends on the magnetic field, chemical potential and temperature. This formula recovers previously studied cases that are solely magnetic field dependent. This renders us to unveil the continuous effect of impurity on the magneto-transport in Weyl semimetals, which is beyond the scope of previous studies. The conclusions are summarized as follows: (i) The magneto-conductivities are sensitive to disorder. \( \sigma_{xx} \) is enhanced by increasing the impurity density; while \( \sigma_{zz} \) is depressed dramatically. (ii) This distinct difference between \( \sigma_{xx} \) and \( \sigma_{zz} \) reveals different scattering mechanisms for the transverse and longitudinal transport. The former behaves like an insulator in which impurity assists the hopping between the neighbors of localized states in the x–y plane, while the latter is conducted like in a metal so that the impurity induces larger resistivity and lower conductivity. (iii) With varying the angle \( \phi \) between \( E \) and \( B \), the magneto-conductivity is a combination of the longitudinal and transverse magneto-conductivities. Our calculated magneto-conductivity recovers the angle-dependent measurements very well. Particularly, we study the effect of impurity on the angular dependence of magneto-conductivity. It is found that the high impurity density suppresses the positive longitudinal magneto-conductivity (i.e., negative longitudinal magneto-resistivity (NLMR)). (iv) As a result of (iii), the NLMR induced by chiral anomaly (around at \( E \parallel B \)) may be buried and cannot be observed when the impurity density is high. We thus comment that heavy doping should be avoided in order to observe the NLMR induced by the chiral anomaly. It is interesting to note that introducing disorder into Weyl semimetals may provide a helpful way to manipulate the valley degree in the realm of valleytronics. This present work may deepen our understanding of the magneto-transport in Weyl semimetals.

We would like to put some words on the perspective of type-II Weyl semimetals.\(^{7,34} \) The Weyl cones are strongly
tilted so that the Weyl nodes exist at the touching points of the electron and hole pockets on the Fermi surface. Compared to the Type-I Weyl semimetal studied in the work, Type-II Weyl semimetals develop anisotropic field dependence due to the tilted Weyl cones with respect to the \( B \) field.\(^{35-37} \) which complicates the structures of Landau levels. At a critical angle between the \( B \) field and the tilt, the Landau level spectrum may collapse.\(^{35} \) Moreover, the structure of Landau levels is quite different when the \( B \) field is parallel to the \( z \) or \( x \) direction.\(^{36} \) This complicated dispersion may induce more complicated consequences on the transport measurements which need complete new treatments rather than the present simple model. What we can speculate is that the dependence of the magneto-transport on the angle of the \( B \)-field is more sensitive and stronger than those in type-I Weyl semimetal in which the model includes one pair of Weyl nodes. Since there should be even a pair of Weyl nodes in type-II Weyl semimetal, the present model should be extended at least to include two pairs. We are sure that the inter-valley scattering should play an important role as well. However, the likely presence of electron pockets and hole pockets on the Fermi surface complicates the speculation.

**SUPPLEMENTARY MATERIAL**

See supplementary material for the derivations of Eqs. (5), (18), (20) and other equations in the main text and the derivations of matrix \( U \), magnetoconductivities and form factors.

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