Aslamazov Larkin Conductivity in Weyl Semimetals

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The Aslamazov-Larkin conductivity (ALC) in the Weyl semi-metals is calculated both for Type-I and Type II for different dimensionality and magnetic fields. The ALC strongly depends on the tilt parameter of the dispersion relation cone. While the 3D and 2D the Aslamazov -Larkin conductivity slightly depends on tilt parameter in Type-I phase and increases in the Type-II phase the 1D ALC decreases in Type-I phase up to zero close to border between Type-I and Type-II phases. Results are discussed in light of the resent experiments on the layered HfTe₅ Weyl semi-metals. It is concluded the one dimensional AL conductivity well explained the experimental data.

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INTRODUCTION.

Dispersion relation near Fermi surface in recently synthesized two and three dimensional Weyl (Dirac) (WSM/DSM) semi-metals [1], [2], [3] is qualitatively distinct from conventional metals, semi - metals or semiconductors in which all the bands are parabolic. In type I Weyl semi-metals (WSM), the band inversion results in Weyl points in low-energy excitations being anisotropic massless ”relativistic” fermions with dispersion relation of these semimetals looks as a tilted Dirac cone. In Type I WSM/DSM the tilt angle of the cone is smaller than some critical value [4] while recently discovered layered transition-metal dichalcogenides is defined as type-II WSM/DSM [5]. In this latter case, the Dirac cone exhibits a strong tilt, so that it can be characterized by a nearly flat band at Fermi surface. The type-II WSM also exhibit exotic topological properties different from the type-I ones, such anti-chiral effect of the chiral Landau level and novel quantum oscillations. Graphene is a prime example of the type I WSM, while materials, like layered organic compound were long suspected to be a 2D type-II Dirac fermions. Several materials were observed to undergo the type-I to II transition while doping or pressure is changed [7], [8], [9]. Theoretically physics of the topological (Lifshitz) phase transitions between the type I to type II Weyl semi-metals were considered in the context of superfluid phase in $^3$He, layered organic materials in 2D and 3D Weyl semi-metals [10], [11], [12]. The pressure modifies the spin orbit coupling that in turn determines the topology of the Fermi surface of these novel materials [13].

Many Weyl and Dirac materials are known to be superconducting. A detailed study of superconductivity in WSM/DSM under hydrostatic pressure revealed a curious dependence of critical temperature of the superconducting transition on pressure. In particular the critical temperature in some of these systems like $HfTe_5$ show a sharp maximum as a function of pressure [14]. This contrasts with generally smooth dependence on pressure in other superconductors (not suspected to be Weyl materials) like cuprates. Magnetic properties in these topological superconductors are also very different from that in conventional superconductors. In particular the Abrikosov parameter used to distinguish between the superconductivity of the first from the second type depends on the cone tilt and may totally change magnetic properties varying from first kind superconductor (like clean metals) to the second kind. The critical fields, coherence lengths magnetic penetration depths and the Ginzburg number characterizing the strength of fluctuations strongly depend on cone tilt [15], [16]. It reveals an extremely important relation between the cone title and fluctuation in WSM/DSM superconductors [17]. We begin with a brief review of studies of fluctuations in superconductors. The subject was initiated in the work of Aslamazov and Larkin [18] calculated the conductivity of fluctuating Cooper pairs in zero magnetic field. Maki and Thomson [19–21] included effects of electron scattering off the fluctuations. It was found that there is another badly divergent contribution known as anomalous Maki-Thomson correction. Physically, this correction is connected with the coherent scattering of the electrons by the impurities and analogous to the weak localization correction. The divergence can be removed by introducing a pair-breaking rate. It should be noted that, the experimental results at temperature close to the critical $T_c$ can be described by the Aslamazov-Larkin term only. This suggests that the pair-breaking rate is relatively large in real superconductors. Later, Thomson and Maki returned to the issue and evaluated fluctuation correction to the normal conductivity in finite fields [22], [23]. A theory of transport phenomena in the fluctuation region in the dirty, clean and super clean limits was developed by Aronov et al. [24]. Their consideration was based on the Ginzburg-Landau equations and, thus, is applicable for relatively small fields. Experimentally the broadened of the normal-to-superconductivity transition curve caused by the superconducting fluctuations was studied in a lot of WSM/DSM materials while the 2D Aslamazov Larkin conductivity was detected in 3D WSM material $Cd_2As_3$ [25]. In the present paper we study of Aslamazov-Larkin conductivity (ALC) in Weyl/Dirac semimetals. The focus generally is on the dependence of the ALC in the cone tilt parameter and consequently on the ALC in the vicinity of the transition from type-I to type-II WSM.

MODEL

Since the Weyl material typically possesses several sublattices the effect of the topological transition on superconductivity was exemplified using the simplest possible model with just two sublattices. The band structure near the Fermi level of a 3D layered Weyl semi-metal is well captured by the non-interacting massless Weyl dispersion relation with the ”in plane” Fermi velocity $v$ (assumed to be isotropic in the $x − y$ plane) and conventional parabolic term on $z$—direction [26], [27], [28], [29] (see Fig.1) and described by the Hamiltonian
FIG. 1. Geometry of the problem

\[ K = \int_r^s \psi^+_\alpha(r) \bar{R}_{\alpha\beta} \psi^\beta_\beta(r) \]  

\[ \bar{R}_{\gamma\delta} = -i\hbar v^\gamma \sigma^\gamma_{\gamma\delta} + \left( -i\hbar w^\gamma - \mu + \frac{p_z^2}{2m_z} \right) \delta_{\gamma\delta}. \]  

Here \( \mu \) is the chemical potential, \( p_z = -i\hbar \nabla_z \), \( \sigma \) are Pauli matrices in the sublattice in the sublattice space in the WSM layers, with just two sublattices denoted by \( \alpha = 1, 2 \) and \( s \) is spin projection. The velocity vector \( w \) defines the tilt of the Dirac dispersion cone. The graphene-like dispersion relation for \( w = 0 \) represents the type I Weyl semi-metal, while for the velocity \( |w| > v \), the material becomes a type II Weyl semi-metal.

We restrict our self to the case of just one left handed and one right handed Dirac points, typically but not always separated in the Brillouin zone. Generalization to include the opposite chirality and several "cones" is straightforward. We assume that different valleys are paired independently and drop the valley indices.

**GINZBURG-LANDAU THEORY.**

The effective electron-electron attraction due to the electron-phonon attraction opposed by Coulomb repulsion (pseudopotential) mechanism creates pairing below \( T_c \). Assuming the singlet s-channel interaction with essentially local interaction in the layers one obtains the set of Gor'kov equations [27] and their microscopic Ginzburg-Landau expansion [28], [29] for arbitrary Dirac cone tilt parameter \( \kappa \) in the form:

\[ F = \int d^3r \left\{ D_0(\mu) f(\kappa) \left( \xi^2(\kappa) |\partial_i \Delta|^2 - \tau(\kappa) |\Delta|^2 + \frac{\beta(\kappa)}{2} |\Delta|^4 \right) + \frac{(\nabla \times A)^2}{8\pi} \right\}, \]  

Where \( \tau(\kappa) = 1 - T/T_c(\kappa) \), \( \partial_i = -i\hbar \nabla_i - 2eA/c, T_c(\kappa) = 1.14\Omega \exp(-1/\lambda_0 f(\kappa)). \)

The index \( i = x, y, z \), \( D_0(\mu) = \frac{\sqrt{2m_0} e^{\mu^2/2}}{12\pi^2 \hbar^2 v^2} \), is the DOS for layered graphene (\( \kappa = 0 \)), dimensionless constant \( \lambda_0 \) is the electron-electron strength for zero tilt parameter \( \kappa \). (We use below units with \( c = \hbar = 1 \)). The tunneling of the electrons moving between the superconducting layers via dielectric streak described by the effective mass \( m_z \) of the electrons moving along the z axis. Within tight binding model the effective mass is estimated as

\[ m_z = \frac{m_e s^2 \exp(d/s)}{d^2} \]  

where \( m_e \) is the mass of free electron, \( d \) is the distance between layers of thickness \( s \), see Fig.1. The GL coefficients in Eq[2] have been calculated for all values of cone tilt parameter \( \kappa \) in our Ref.[29]. In particular for the type I WSM,
\[ f (\kappa) = \frac{1}{(1 - \kappa^2)^{3/2}}. \tag{4} \]

In the type II phase, \( \kappa > 1 \), the Fermi surface becomes open, extending over the Brillouin zone, and the corresponding expression is:

\[ f = \frac{\kappa^2}{\pi (\kappa^2 - 1)^{3/2}} \left\{ 2 \sqrt{1 + \kappa} - 1 + \log \left[ \frac{2 (\kappa^2 - 1)}{\kappa (1 + \sqrt{1 + \kappa})} \epsilon \right] \right\}. \tag{5} \]

Here \( \epsilon \) is an ultraviolet cut off parameter \( \epsilon = a \Omega / \pi w \), where \( a \) is an interatomic spacing, \( \Omega \) is the phonon frequency around the Fermi energy.

The critical temperature \( T_c (\kappa) \) has the sharp spike at the border between Type-I and Type-II states at \( \kappa \rightarrow 1 \) where the topology of the Fermi surface undergoes the Lifshitz 2.5 type transition \[28\].

**ASLAMAZOV-LARKIN CONDUCTIVITY.**

Using the standard time dependent Ginzburg Landau equations \[30\] with the external \( \delta \) correlated random force responsible for thermal fluctuations:

\[
\Gamma \left( \frac{\partial}{\partial t} + 2i e \Phi \right) \Delta = D_0 (\mu) f (\kappa) \left( \xi_i (\kappa) \xi_j (\kappa) \partial_i \partial_j - \tau (\kappa) + \beta (\kappa) |\Delta|^2 \right) \Delta + \Gamma T \zeta (r, t) \tag{6}
\]

\[
J = 2ei D_0 (\mu) f (\kappa) \xi_i^2 (\kappa) \partial \Delta^* (r) + c.c. + \sigma_n \left( - \frac{\partial A}{\partial t} - \nabla \Phi \right). \tag{7}
\]

where \( \Gamma = D_0 (\mu) f (\kappa) \pi / 8 T_c (\kappa) \) is the relaxation constant \( \Phi \) is the electric potential, \( A \) is the vector potential, \( \sigma_n \) is the normal conductivity.

Taking order parameter in the form \( \Delta (r, t) = \Delta_0 (r) + \Delta_1 (r, t) \) where \( \Delta_0 (r) \) is the order parameter in the absence of the electric field, and substituting it into Eqs.\[6,7\] one obtains in the case of zero magnetic field for the dc superconducting current:

\[
J_\alpha = 4 \Gamma e^2 k_\alpha^2 \xi_\alpha^4 (\kappa) \int k \frac{k_\alpha}{\xi_\alpha^2 (\kappa) k_\beta^2 + \tau} E_\beta \frac{\partial}{\partial k_\beta} \left\langle |\Delta_0 (k)|^2 \right\rangle \tag{8}
\]

where

\[
|\Delta_0 (k, \kappa)|^2 = \frac{T_c (\kappa)}{D_0 (\mu) f (\kappa) (\xi_\alpha^2 (\kappa) k_\alpha^2 + \tau (\kappa))} \tag{9}
\]

Using \[8\] one obtains for the Aslamazov-Larkin conductivity:

\[
\sigma_{\alpha \alpha} = e^2 \pi \xi_\alpha^4 (\kappa) \int k \frac{k_\alpha^2}{(\xi_\alpha^2 (\kappa) k_\alpha^2 + \tau (\kappa))} \tag{10}
\]

Performing integration one obtains for different dimensionality. For 3D case the tilt and temperature dependence of the ALC has the form:

\[
\sigma_{\alpha \alpha} (\kappa) = \frac{e^2 \xi_\alpha (\kappa)}{64 \pi \xi_\beta (\kappa) \xi_\gamma (\kappa) \tau (\kappa)^{1/2}}, \sigma_{zz} (\kappa) = \frac{e^2 \xi_z (\kappa)}{32 \xi_x (\kappa) \xi_y (\kappa) \tau (\kappa)^{1/2}} \tag{11}
\]
FIG. 2. ALC of the 3D Weyl superconductors as a function of the tilt cone parameter $\kappa$

Where $\alpha = x, y$.

The reduced ALC $\sigma_{xx} = \sigma_{xx}(\kappa) / \sigma_{xx}(0)$ is shown in Fig. 2 and demonstrates strong dependence on the tilt parameter in Type-II phase.

For 2D and 1D cases one obtains from Eq. 10:

$$\sigma^{2D}_{\alpha\alpha}(\kappa) = \frac{e^2 \xi_{\alpha}(\kappa)}{16\pi d \xi_\beta(\kappa) \tau(\kappa)}; \sigma^{1D}_{\alpha\alpha}(\kappa) = \frac{\pi e^2 \xi_{\alpha}(\kappa)}{16S \tau(\kappa)^{3/2}}$$

(12)

and presented in Figs. 3, 4

Here $d$ is the film thickness and $S$ is the cross-section of the wire, $\alpha \neq \beta$.

FIG. 3. ALC of the 2D Weyl superconductors as a function of the tilt cone parameter $\kappa$

ASLAMAZOV-LARKIN CONDUCTIVITY IN MAGNETIC FIELD.

The ALC in magnetic field was calculated earlier for both for isotropic system [31] and for layered structure with Josephson interaction between the layers [32] but these results cannot be used in WSM with a strong anisotropies of
the coherence lengths. In our case with the external magnetic field directed parallel to the $z$–axis and the vector-potential $A = (Hy,0,0)$, the ac ALC can be calculated from the dissipative-fluctuation theorem in the Kubo form:

$$\sigma(\omega,\kappa) = \frac{2\pi}{\omega} \tanh \left( \frac{\omega}{2T} \right) \int dy \left\langle J_\omega (0, \frac{y}{2}, 0) J_{-\omega} (0, -\frac{y}{2}, 0) \right\rangle$$

(13)

Here $J_\omega (k, \frac{y}{2}, k')$ is the Fourier transform of the current density with the fluctuating order parameter in the form:

$$\Delta (\mathbf{r}, t) = \int dk dp \sum_n \Delta_n,p,k (t) \phi_n \left( y - \frac{c}{2eH} k_x \right) e^{-ipz - ikx}$$

(14)

where $\phi_{n,y_0} (y)$ are the normalized eigenfunction of the electron in magnetic field

$$\phi_{n,y_0} (y) = \frac{\exp \left( -(y - y_0)^2 / 2a_H^2 \right)}{\pi^{1/4} \sqrt{2^n a_H n!}} H_n \left( \frac{y - y_0}{a_H} \right);$$

(15)

with the eigenvalues

$$\varepsilon_{n,k_z} = \omega_H (2n + 1) + \xi_x^2 k_z^2, \omega_H = \frac{eH \sqrt{2}}{c} \xi_x (\kappa) \xi_y (\kappa).$$

(16)

Here $H_n$ is the Hermite polynomial, $a_H^2 = c \xi_y / \sqrt{2eH} \xi_x$.

Substituting Eq. (14) into Eq. (7) one obtains for the current density after Fourier transform:

$$J_\omega \left(0, \frac{y}{2}, 0\right) = -2ieD (\kappa) \xi_y^2 \Pi;$$

(17)

$$\Pi = \sum \Pi_{nm};$$

(18)

$$\Pi_{nm} = \int_{-\infty}^{\infty} dk dp d\omega \left[ \Delta_{n,p,k} \left( \omega + \frac{\Omega}{2} \right) \Delta^*_{m,p,k} \left( \omega - \frac{\Omega}{2} \right) F_{n,m} \left( \frac{y}{2}, k, p \right) \right]$$

(19)

where
\[ F_{n,m}(y,k_1,k_2) = \phi_n \left( y - \frac{c}{2eH}k_1 \right) \phi'_m \left( y - \frac{c}{2eH}k_2 \right) - \phi_m \left( y - \frac{c}{2eH}k_2 \right) \phi'_n \left( y - \frac{c}{2eH}k_1 \right) \]  

(20)

Using the random force correlator in the right side of the Eq.(6) one obtains for the correlator of the order parameter

\[ |\Delta_n|^2 = T \left[ D(\kappa) \left( \omega^2 + E_{n,p,k}^2 \right) \right]^{-1} \], where \( E_{n,p,k} = \varepsilon_{n,p,k} + \tau \).

Substituting this correlator into the Kubo relation for conductivity and performing the integrations one obtains for ALC in magnetic field

\[ \sigma_{yy}(\kappa, 0) = \frac{e^2 \xi_y(\kappa)}{16\pi \xi_x(\kappa) \xi_z(\kappa)} \sum (n + 1) \left( \frac{\sqrt{\omega H (2n+3) + \tau(\kappa)}}{\sqrt{\omega H (2n+1) + \tau(\kappa)}} \right) \left( \frac{1}{2\sqrt{\omega H (2n+2) + \tau(\kappa)}} \right) \]  

(21)

It gives in the limit of a small magnetic fields \( H << H_{c2} \), where \( H_{c2} = \Phi_0 \tau(\kappa)/2\pi \xi_x(\kappa) \xi_z(\kappa) \) is the upper critical magnetic field, \( \Phi_0 \)is the unit flux

\[ \sigma_{yy}(\kappa, 0) = \frac{e^2 \xi_y(\kappa)}{64\pi \xi_x(\kappa) \xi_z(\kappa)} \frac{1}{\tau^{1/2}(\kappa)} \left\{ 1 - \frac{23}{16} \left( \frac{H}{H_{c2}} \right)^2 \right\} \]  

(22)

In a strong magnetic field ALC diverges close to the \( H_{c2} \). In this case the main term in the sum is the second ones with \( n = 0 \)

\[ \sigma_{yy}(\kappa, 0) = \frac{e^2 \xi_y(\kappa)}{16\pi \xi_x(\kappa) \xi_z(\kappa)} \sqrt{\frac{H_{c2}^0}{H - H_{c2}}} \]  

(23)

and diverges at the coexistence line.

**DISCUSSION AND CONCLUSIONS.**

The Aslamazov-Larkin conductivity in the Weyl metals is described by the Eqs. [11][12] in the absence of the magnetic field and [21][22][23] in the presence of the magnetic field for different dimensionality and demonstrate typical for AL conductivity temperature and field dependences. However, the coherence lengths are the cone slope dependent (see Figs.2-4 and Fig. 6) demonstrated the main results. While for 3D and 2D the Aslamazov-Larkin conductivity demonstrate similar dependence on the tilt parameter (weak variation in Type-I phase and essential increase in the Figs.2-4 and Fig. 6) demonstrated the main results. While for 3D and 2D the Aslamazov-Larkin conductivity demonstrate similar dependence on the tilt parameter (weak variation in Type-I phase and essential increase in the Type-II phase, the ALC in 1D looks very different significantly reduces close to the border between Type-I and Type-II phases.

As it was predicted in Ref. [7], [8], [9] the static pressure can control the cone slope parameter \( \kappa \) in Weyl semimetals. In particular, the type-I WSM realized at large range of the pressures, while the type-II appears at low pressures. Recently, the resistance in the WSM was measured in the 3D layered material \( HfTe_5 \) [14].

The crystal structure of \( HfTe_5 \) has been determined by powder X-ray diffraction experiments [13]. Trigonal prismatic chains of \( HfTe_5 \) run along the a axis, and these prismatic chains are linked via parallel zigzag chains of Te atoms along the c axis to form a 2D sheet of \( HfTe_5 \) in the ac plane (along the z direction in our notation). The sheets of \( HfTe_5 \) stack along the b axis, forming a layered structure [14] where metallic atoms are surrounded by dielectric as it is shown in Fig. 5 ([14]).

This material shows a well pronounced dependence of the resistivity on the static pressure. In particular, it was established that below 4.7 GPa the type-I WSM is realized while the type-II is realized above 6.1 GPa. In the range 4.7 - 6.1 GPa, the type-II and type-I Dirac cones coexist. The critical temperature non-monotonic depending on the pressure allows to assume that the pressure is directly related to the cone slope \( \kappa \) parameter. It is naturally to assume that the Type-II WSM is established at large pressure and small \( \kappa \).
FIG. 5. Crystal structure of the $\text{HfTe}_5$ in (100) projection. Black balls mark metal atoms, while the rest non-metals. The zig-zag-like layers run along the $b$-axis ($z$ axis in our notations).

FIG. 6. The experimental results (points) and the theoretical fitting by the (solid curve) for AL conductivity demonstrates the 1D character of the fluctuations inside the $\text{HfTe}_5$ layers.

The resistivity was measured in the direction perpendicular to the layers. The comparison with the theory is presented in Fig.6 and demonstrates a good agreement with 1D ALC formulae. It shows that the conductivity obeyed the 1D AL Eq. 12 (see Fig.4). Measurements of the magnetoresistance when the current flows perpendicular to the layers (direction) show that its temperature dependence is well described by the one-dimensional formula for Aslamazov-Larkin fluctuations. This indicates that the current flows along the chains of the metal atoms forming one-dimensional channels crossing the layers.
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