Andreev Reflection in Weyl Semimetals

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We theoretically study low energy electric transport in a junction consisting of a Weyl semimetal and a metallic superconductor. The characteristic features of the differential conductance depend on the relative directions between the current and the vector connecting the two Weyl points. When the electric current is perpendicular to the vector, the conductance spectra are sensitive to the direction and the amplitude of magnetic moment at the junction interface. This is a direct consequence of the chiral spin configuration on the fermi surface near the Weyl points.

1. Introduction

Weyl semimetals are a novel topological material in three-dimension.¹⁻⁷ In the bulk, Weyl semimetal has the gapless spectra characterized by even number of Weyl points with opposite topological charges in Brillouin zone, which leads to unusual transport properties.⁸⁻¹⁵ Topological properties in these cases result from the separation of the individual Weyl points in the absence of either time-reversal or inversion symmetry. Various theoretical model has been discussed to realize the semimetal phase.¹⁶⁻¹⁸ In experiments, the semimetal is realized in multilayers of GeTe/Sb₂Te₃¹⁹,₂⁰ and Cd₃As₂.²¹⁻²³ Recently a possibility of superconductivity has been discussed in several theoretical studies.²⁴⁻²⁶ Unconventional Cooper pairing symmetry would be expected because of the chiral spin structure on the fermi surface. Although finding the superconductivity within existing semimetals might be difficult, realizing the superconducting correlation there is basically possible by injecting Cooper pairs due to the proximity effect. As the first step in this research direction, we address the Andreev reflection in Weyl semimetal in the present paper.

In this paper, we theoretically study the low energy transport through a junction consisting a Weyl semimetal and a metallic superconductor. The differential conductance of the Weyl-semimetal/superconductor junction is calculated from the normal and the Andreev reflection coefficients of the junction. We consider a 2×2 simple Hamiltonian which describes the electronic structure of a Weyl semimetal breaking the time-reversal symmetry. Within this model, the semimetallic excitation is characterized by the linear energy-momentum dispersion relations near the two separated Weyl points. The spin structures at the two fermi surfaces are characterized by the opposite spin chiral texture to each other. When the two Weyl points stay at ±K₀ with K₀ = (0, 0, k₀), the topological surface states appears on the four surfaces parallel to K₀ direction. We consider two types of junction: the current parallel to K₀ and the current perpendicular to K₀. When the potential barrier at the interface is spin-independent, the conductance spectra in the junctions are similar to those in the usual normal-metal/superconductor junctions. When the potential at the interface is spin active, on the other hand, the conductance spectra depend sensitively on types of the junction and directions of the magnetic moment. In particular, for the cases with magnetic moment K₀, both the Andreev reflections are suppressed due to the chiral spin structure on the fermi surface. The spin-flip potentials at the interface relax the spin mismatch in the reflection process.

2. Weyl Semimetal

To describe the electronic states in Weyl semimetals, we use a simple model given by

\[ H_W = \sum_{\alpha, \beta} \int d\mathbf{r} \psi_\alpha^\dagger(r) \left[ -\frac{\hbar^2}{2m_W} \nabla^2 + k_0^2 \right] \hat{\sigma}_z - i\lambda(\partial_\alpha \hat{\sigma}_x + \partial_\beta \hat{\sigma}_y) - \mu_W \hat{\sigma}_\alpha \right] \psi_\beta(r), \]

where \( \psi_\alpha^\dagger(r) \) (\( \psi_\beta(r) \)) is the creation (annihilation) operator of an electron with spin \( \alpha \) at \( \mathbf{r} \), \( \nabla \) is the three-dimensional Laplacian, \( m_W \) is the effective mass of an electron, \( \lambda \) denotes the coupling constant of the spin-orbit interaction, and \( \mu_W \) is the chemical potential measured from the Weyl point. We originally begin with the spin-degenerate two-band model as shown in Fig. 1(a). The Zeeman field decreases (increases) the energy of the spin-up (spin-down) band. Large enough Zeeman fields result in the inverted band structure. The effects of the Zeeman field is taken into account through \( k_0 \) in Eq. (1). The Pauli’s matrices \( \hat{\sigma}_j \) for \( j = x, y, z \) represent the real spins of an electron. The unit matrix in spin space is \( \hat{\sigma}_0 \). By neglecting the two bands away from the fermi level, the electric structure are described by Eq. (1). In the Fourier representation, Eq. (1) becomes

\[ H_W(k) = \begin{bmatrix} \epsilon_k - \mu_W & \lambda(k_x - ik_y) \\ \lambda(k_x + ik_y) & -\epsilon_k - \mu_W \end{bmatrix}, \]

with \( \epsilon_k = (\hbar^2/2m_W)(k^2 - k_0^2) \). The energy dispersion and the wave functions are obtained as

\[ \begin{pmatrix} \alpha_k \\ \beta_k(k_x + ik_y)/p \end{pmatrix}, \quad \begin{pmatrix} -\beta_k(k_x - ik_y)/p \\ \alpha_k \end{pmatrix}, \]

for \( E_W^W - \mu_W \) and \(-E_W^W - \mu_W \), respectively. Here we define following quantities,

\[ E_k^W = \sqrt{\epsilon_k^2 + (\lambda p)^2}, \quad p = (k_x, k_y, 0), \]

\[ \alpha_k = \sqrt{2} \left( 1 + \frac{\epsilon_k}{E_k^W} \right), \quad \beta_k = \frac{1}{\sqrt{2}} \left( 1 - \frac{\epsilon_k}{E_k^W} \right). \]
The spin expectation values are calculated to be $\langle \sigma_s(x) \rangle = \frac{\hbar k_{s(x)}^2}{E_k^W} + i \frac{\hbar k_{s(x)}}{E_k^W}$, in units of current parallel to the \( \lambda \) axis indicated by a broken line. (b) We consider two types of junction: current parallel to the \( z \) axis and that parallel to the \( x \) axis. The topological surface states appear on the surfaces perpendicular to the \( x \) axis and those to the \( y \) axis.

The two Weyl points appear at \( \pm K_0 \). When we consider low energy transport around the chemical potential, the energy band with \( E_k^W = \mu_W \) carries the electric current. The group velocity of Weyl semimetal is anisotropic. The velocities are represented by

$$v_{s(x)} = \frac{\hbar k_{s(x)}}{m_W E_k^W} \frac{e_k}{E_k^W}, \quad v_z = \frac{\hbar k_z}{m_W E_k^W},$$

(6)

in the \( s(x) \) and the \( z \) direction, respectively. In addition, the expectation value of spin \( S = (h/2) \hat{\sigma} \) are calculated to be

$$\langle S_{s(x)} \rangle = \frac{\hbar k_{s(x)}}{E_k^W}, \quad \langle S_z \rangle = \frac{\hbar k_z}{E_k^W},$$

(7)

in units of \( h/2 \).

The Hamiltonian in the hole space is represented by $-H^W_k(-k)$. The wave function at $E = -E_k + \mu_W$ is given by

$$\begin{pmatrix}
\alpha_k \\
-\beta_k(k_x - ik_y) / 1
\end{pmatrix}.$$  

(8)

The spin expectation values are calculated to be

$$\langle S_{s(x)} \rangle = \frac{\hbar k_{s(x)}}{E_k^W} = \frac{\hbar k_z}{E_k^W},$$

(9)

in the hole space.

### 2.1 Fermi surface

To study transport properties unique to the Weyl semimetals, we need to set the chemical potential $\mu_W$ to be small enough values. We show the shape of the fermi surface at $\lambda = 0.5\epsilon_0$ for several $\mu_W$ in Fig. 2, where the line connects the equal energy points in the Brillouin zone. For $\mu_W/\epsilon_0 < 0.4$, the two disconnected fermi surfaces enclose the two Weyl points at $(0, 0, \pm k_0)$. We only show the fermi surface around $K_0 = (0, 0, k_0)$ in Fig. 2. At $\mu_W/\epsilon_0 = 0.2$, the shape of fermi surface is still distorted. The fermi surface becomes more ellipsoidal for smaller $\mu_W$. In this paper, we fix the parameters as $\lambda = 0.5\epsilon_0$ and $\mu_W = 0.1\epsilon_0$.

### 2.2 Wave function at a fixed energy

It is possible to consider two types of junction: (i) the current in the $z$ direction and (ii) the current in the $x$ direction. We first discuss the wave function for the current in the $z$ direction. At an energy $E > 0$, such wave function in the $z$ direction proportional to $e^{i(k_x x + k_y y)}$ is described by

$$\Psi_{w(x)}(z \leq 0) = \begin{pmatrix} f_{e} & 0 \\ 0 & 0 \end{pmatrix} \left(e^{ik_x z} a^+ + e^{-ik_x z} A^+\right)$$

(10)

$$+ \begin{pmatrix} \gamma_+ & 0 \\ 0 & 0 \end{pmatrix} \left(e^{-ik_x z} a^- + e^{ik_x z} A^-\right)$$

$$+ \begin{pmatrix} 0 & f_{e}^* \\ 0 & 0 \end{pmatrix} \left(e^{ik_x z} b^+ + e^{-ik_x z} B^+\right)$$

$$+ \begin{pmatrix} 0 & 0 \\ -\gamma_+ & f_{e}^* \end{pmatrix} \left(e^{ik_x z} b^- + e^{-ik_x z} B^-\right),$$

(11)

$$f_{e(h)} = \mu_W + (-\mu_W + E) \sqrt{D_{e(h)}^w}, \quad \gamma_+ = \lambda(k_x \pm ik_y),$$

(12)

$$k_{e(h)}^z = \sqrt{k_0^2 - p^2 \pm (2m_W h^2)^{1/2}} \left|D_{e(h)}^w\right|,$$

(13)

where $c_{e(h)}$ and $c_{e(h)}^*$ are factors which normalize the wave functions. The coefficients $a^+$ ($b^+$) are the amplitudes of wave function incoming into the junction interface as an electron (hole). While $A^+$ ($B^+$) are the amplitudes of wave function outgoing from the junction interface as an electron (hole). We note that $E_k^W$ was replaced by $\mu_W + E$ in the electron space. At the same time, $\epsilon_k$ was replaced by $\pm D_{e(h)}^W$ for $k_{e(h)}^z$ channel. In the hole space, we have applied relations $E_k^W \rightarrow \mu_W - E$ and $\epsilon_k \rightarrow \pm D_{e(h)}^W$ for $k_{e(h)}^z$ channel. In this way, we obtain Eq. (10).

When the current flows in the $x$ direction, on the other hand, the wave function in the $x$ direction proportional to $e^{i(k_x x + k_y y)}$
In the hole space, we have applied Eq. (14). In the electron space, the wave function of the bound states are the eigen states of \( \hat{\Sigma} \) by using the wave function in the complex wave number \( \lambda \) under the dispersion of the surface bound states as

\[
\begin{align*}
\Psi_{w}(x \leq 0) &= \frac{1}{\tilde{z}_{k_{y}}} \begin{pmatrix} f_{e_{-}}^{+} & e^{-ik_{y}x}a^{+} & 1 \end{pmatrix} e^{ik_{y}x}a^{-}, \\
+ \frac{1}{\tilde{z}_{-k_{y}}} \begin{pmatrix} f_{e_{+}}^{+} & \gamma_{-k_{y}+}^{*} & e^{-ik_{y}x}A^{+} \end{pmatrix} e^{ik_{y}x}A^{-}, \\
+ \frac{1}{\tilde{z}_{k_{y}}} \begin{pmatrix} f_{h_{+}}^{+} & e^{-ik_{y}x}b^{+} & 1 \end{pmatrix} e^{ik_{y}x}b^{-}, \\
+ \frac{1}{\tilde{z}_{-k_{y}}} \begin{pmatrix} f_{h_{-}}^{+} & \gamma_{-k_{y}+}^{*} & e^{-ik_{y}x}B^{+} \end{pmatrix} e^{ik_{y}x}B^{-},
\end{align*}
\] (20)

The topologically protected surface bound states appear at the parallel directions to the interface are conserved in the above conditions. In addition to this, we need to add the mass of an electron \( (\text{the chemical potential}) \) in the superconductor. The wave function in the \( z \) direction, for example, is described by

\[
\Psi_{w}(z = 0) = \Psi_{S}(z = 0),
\] (27)

\[
\begin{pmatrix} e^{i\sigma_{z}} & 0 \\
0 & e^{i\sigma_{z}} & -\Delta_i\hat{\sigma}_{y} \\
0 & -\Delta_i\hat{\sigma}_{y} & -\epsilon_{-k_{y}}\hat{\sigma}_{x} & 0 \\
-\Delta_i\hat{\sigma}_{y} & -\Delta_i\hat{\sigma}_{y} & -\epsilon_{-k_{y}}\hat{\sigma}_{x} & 0 \\
\end{pmatrix}
\]

\[
\begin{pmatrix} \xi_{k}^{2}k^{2} - m_{s}^{2} \end{pmatrix}
\] (23)

\[
\begin{pmatrix} 1 + \Omega & 0 \\
0 & 1 - \Omega \\
\end{pmatrix}
\]

\[
\begin{pmatrix} u_{0} & 0 \\
0 & -u_{0} \\
\end{pmatrix}
\]

\[
\begin{pmatrix} D_{\hat{z}} & \epsilon_{e}^{(i)} \hat{\sigma}_{y} \\
\epsilon_{e}^{(i)} \hat{\sigma}_{y} & D_{\hat{z}} \end{pmatrix}
\]

\[
\begin{pmatrix} 0 \\
0 \\
-D_{\hat{z}} \hat{\sigma}_{y} \\
D_{\hat{z}} \hat{\sigma}_{y} \\
\end{pmatrix}
\]

\[
\begin{pmatrix} A^{+} \\
A^{-} \\
B^{+} \\
B^{-} \\
\end{pmatrix}
\]

\[
\begin{pmatrix} \hat{r}_{ee} & \hat{r}_{eh} \\
\hat{r}_{eh} & \hat{r}_{hh} \\
\end{pmatrix}
\]

\[
\begin{pmatrix} \hat{a}^{+} \\
\hat{a}^{-} \\
\hat{b}^{+} \\
\hat{b}^{-} \\
\end{pmatrix}
\]

The differential conductance of the junction is calculated...
Such behaviors are well known in the conductance spectra
of the pair potential is fixed at $\mu_\text{S} = 2\varepsilon_0$, $\mu_\text{W} = 0.1\varepsilon_0$, and $\mu_\text{WS} = 0.1\varepsilon_0$.

Fig. 3. The differential conductance of Weyl-semimetal/superconductor
junction is plotted as a function of the bias voltage for the current parallel
to the $z$ axis in (a) and for the current parallel to the $x$ axis in (b). Here we
choose $m_\text{WS} = \mu_\text{S} = 2\varepsilon_0$, $\lambda = 0.5\varepsilon_0$, and $\mu_\text{WS} = 0.1\varepsilon_0$.

$$G_{\text{WS}}(eV) = \frac{e^2}{h} \sum_P \text{Tr} \left[ 1_{ee} - R_{ee} R_{ee}^\dagger + R_{he} R_{he}^\dagger \right]_{E=eV},$$

$$R_{ee} = \left( \begin{array}{cc} \sqrt{v_e^+} & 0 \\ 0 & \sqrt{v_e^-} \end{array} \right) \left( \begin{array}{cc} 1/\sqrt{v_e^+} & 0 \\ 0 & 1/\sqrt{v_e^-} \end{array} \right),$$

$$R_{he} = \left( \begin{array}{cc} \sqrt{v_h^+} & 0 \\ 0 & \sqrt{v_h^-} \end{array} \right) \left( \begin{array}{cc} 1/\sqrt{v_h^+} & 0 \\ 0 & 1/\sqrt{v_h^-} \end{array} \right),$$

$$1_{ee} = \left( \begin{array}{cc} s_e^+ & 0 \\ 0 & s_e^- \end{array} \right),$$

where $v_e^+ (v_h^+)$ is the velocity of $k_e^\pm (k_h^\pm)$ channel at the electron (hole) space in the Weyl semimetal and $P$ denotes the
momenta parallel to the interface. We define that $s_e^+$ is 1 for the
propagating $k_e^+$ channel and 0 for the evanescent $k_e^-$ channel
in the electron space.

3. Differential Conductance

Throughout this paper, we fix material parameters in the superconductor as $m_\text{S} = m_\text{W}$ and $\mu_\text{S} = 2\varepsilon_0$. These parameters only modify the transmission probability of the junction. The amplitude of the pair potential is fixed at $\Delta = 0.01\varepsilon_0$ which gives the smallest energy scale in our model. As discussed in Sec. 2.1, we choose $\mu_\text{W} = 0.1\varepsilon_0$ and $\lambda = 0.5\varepsilon_0$ in the Weyl semimetal. The barrier potentials are parameterized as $z_0 = (m_\text{W}/\hbar^2k_0) V_0$ and $M = (m_\text{W}/\hbar^2k_0) V$.

We first discuss the differential conductance without magnetic barrier at the interface (i.e., $M = 0$). Fig. 2 shows the
differential conductance of Weyl-semimetal/superconductor junction for the current parallel to the $z$ axis in (a) and that
for the current parallel to the $x$ axis in (b). The results are normaliz-
ed to the conductance of Weyl-semimetal/normal-metal junction ($G_{\text{WN}}$) at $eV = 0$. All the results in (a) and (b) show
the dip structure below the gap because the transmission probability
of the junction is less than unity even at $z_0 = 0$ reflecting the
difference in the band structure between the Weyl semimetal and the superconductor. The transmission probability $T_N$ in the normal state is 0.76, 0.48, 0.16 and 0.076 for $z_0 = 0, 1, 3$ and 5, respectively in (a). In (b), $T_N$ is 0.71, 0.31 and 0.05 for $z_0 = 0, 1$ and 3, respectively. The gap structure becomes clearer when we decrease $T_N$ by increasing $z_0$.

4. Spin configuration on the fermi surface within $xz$ plane in spin space. In the junction for the current parallel to $z$, an electron is incident from the two fermi surfaces near $(0, 0, \pm k_0)$ having the opposite spin chiral texture to each other as shown in (a). For the current parallel to $x$ in (b), an electron is reflected into the same fermi surface because $k_x$ and $k_z$ are conserved due to the translational symmetry in $yz$ plane. We only show the spin configuration on the fermi surface around $(0, 0, k_0)$ in (b). The spin is conserved in the normal
device and it becomes opposite direction in the Andreev reflection. The +/- in the bottom of figures indicates the sign of the velocity in the current direction.

$$S(k_x, k_y, k_z^\pm) = \left( \begin{array}{cc} \frac{\Delta k_x}{\mu_\text{W} + E} & \frac{\Delta k_y}{\mu_\text{W} + E} \\ \frac{\Delta k_x}{\mu_\text{W} + E} & \frac{\sqrt{D_z^\pm}}{\mu_\text{W} + E} \end{array} \right),$$

for an incoming electron at $k_x = k_z^\pm$. The wave numbers in the $z$ direction for outgoing channel in the electron space are $k_z^+$ and $-k_z^-$. The former is in the same fermi surface as that in the incident wave, whereas the latter belongs to the opposite fermi surface. In such outgoing channels, $S_x$ and $S_y$ remain unchanged from those in the incoming one. However, $S_z = \sqrt{D_z^+}/(\mu_\text{W} + E)$ and $-\sqrt{D_z^-}/(\mu_\text{W} + E)$ for $-k_z^-$ and $k_z^+$, respectively. In Fig. 4(a), the direction of spin within $xz$ spin plane is illustrated by arrows on the fermi surface at $k_z = 0$. It is possible to obtain the same spin configuration within $yz$ plane at $k_z = 0$. Due to the spin mismatch in $S_z$, the normal reflection to $k_z^+$ is basically suppressed. To conserve the spin direction, therefore, an electron incoming at $k_z^+$ is reflected into $-k_z^-$ in the normal reflection. In the outgoing channels in
On the other hand, they are almost opposite to those in the incoming electron. However, \( S_y \) component depends on the outgoing channels. The Andreev reflection flips the spin of an incoming electron because we assume the spin-singlet superconductor. Thus the Andreev reflection is possible at \(-k^+_n\) belonging to the opposite fermi surface. An incident electron from one fermi surface is reflected into the opposite fermi surface without suffering the spin mismatch. In the junction for the current parallel to \( z \) direction, therefore, the reflection processes are as usual as those in the normal metal. We note that there is no spin mismatch in the \( S_y \) component.

The situation in the junction for the current parallel to the \( x \) axis is more complicated. The spin of an incoming electron at \( k_x = k^+_n > 0 \) channel is

\[
S(k^+_n, k_y, k_z) = \frac{1}{\mu_w + E} \left( \lambda k^+_n, \lambda k_y, -\frac{\lambda k_0}{2} + \sqrt{D^+_h} \right). \tag{36}
\]

On the other hand, they are

\[
S(-k^+_n, k_y, k_z) = \frac{1}{\mu_w + E} \left( -\lambda k^+_n, \lambda k_y, -\frac{\lambda k_0}{2} + \sqrt{D^+_h} \right), \tag{37}
\]

in the normal reflection at \( k_x = -k^+_n \) in the electron space, and

\[
S(k^+_n, k_y, k_z) = \frac{1}{\mu_w - E} \left( -\lambda k^+_n, -\lambda k_y, -\frac{\lambda k_0}{2} + \sqrt{D^-_h} \right). \tag{38}
\]

in the Andreev reflection at \( k_x = k^+_n > 0 \) in the hole space. Although the \( S_y \) component always satisfies the spin selection rule, \( S_x \) and \( S_y \) components violate the selection rule depending on the incident angle. In Fig. 4(b), the spin configuration within \( xz \) spin plane is illustrated by arrows on the fermi surface. We only show the spin structure on the fermi surface around \( k_x = k_0 \) because the wave number in the \( yz \) plane is conserved in the reflection process. When the incident electron has a wave number as indicated by \( A \) in Fig. 4(b), the normal reflection to \( A_1 \) is suppressed due to the spin mismatch in \( S_x \) component but the Andreev reflection to \( A_2 \) is possible. On the other hand, when the incident electron has a wave number as indicated by \( B \), the normal reflection to \( C \) itself is possible but the Andreev reflection to \( C_2 \) is suppressed because of the spin mismatch in \( S_y \) component. At the intermediate incident wave number as indicated by \( B \), both the normal and the Andreev reflections are allowed. Therefore the reflection property depends strongly on the incoming wave number. The results in Fig. 3(b) indicate that the Andreev reflection probability is small near the zero-bias of the spin mismatch in the reflection process. The conductance spectra, however, are expected to be sensitive to spin active potential at the interface.

In Fig. 5, we show the conductance spectra in the presence of the spin dependent potential at the interface, where \( z_0 = 0 \) and \( |M| = 0.1 \). In (a), results for the current in the \( z \) direction is plotted for three directions of \( M \). The conductance spectra are almost unchanged from the results in Fig. 3(a) with \( z_0 = 0 \) because there is no spin mismatch in the reflection processes. However, the conductance spectra for the current parallel to \( x \) depends sensitively on the direction of the magnetic moment \( M \). For example, \( +z \) and \( -z \) mean the magnetic moment points \( +z \) and \( -z \) direction in spin space, respectively. The conductance for \( +x \) and that for \( -x \) are identical to each other. The magnetic moment at the interface drastically modifies subgap spectra for the current parallel to \( x \) direction because it relaxes the spin mismatch in both the normal and the Andreev reflection processes.

The conductance spectra also depends on the amplitude of the magnetic moment. In Fig. 6, we show the conductance spectra for \( |M| = 0.5 \). The peak at \( eV = \Delta \) is suppressed by the magnetic moment for the current parallel to \( z \) direction as shown in (a). The subgap spectra are totally smooth function of \( eV \). In (b), on the other hand, the results show a rich variety of the subgap spectra depending on the directions of magnetic moment. In particular, the amplitudes of the conductance around the zero-bias for \( \pm z \) and \( x \) are larger than those at \( M = 0 \) in Fig. 3(b). This means the large amplitude of the Andreev reflection probability.
4. Discussion

The Andreev reflection means the penetration of a Cooper pair into the Weyl semimetal. When the current is parallel to \( z \) direction, the results in Fig. 3(a) are qualitatively the same as those in normal-metal/superconductor junction as we discussed in Sec. 3. Thus a spin-singlet \( s \)-wave Cooper pair would be dominant in the semimetal. When the current is parallel to \( x \), on the other hand, the spin-flip scattering assists the Andreev reflection as shown in Figs. 5(b) and 6(b). Thus the spin-triplet Cooper pairs are expected in the Weyl semimetal. In addition to this, the orbital part of a Cooper pair is also modified by the reflection. The junction interface can mix the even-parity and odd-parity components because it breaks the translational invariance. Thus the odd-parity spin-triplet component is expected as well as the even-parity spin-singlet one. Moreover a Cooper pair with odd-frequency symmetry might stay in the Weyl semimetal in the dirty limit.\textsuperscript{29,30} To resolve the pairing symmetry, we need calculate the anomalous Green function and analyze it. This would be an interesting issue in the future.

5. Conclusion

We have theoretically studied the differential conductance in the junction of Weyl-semimetal and metallic superconductor. The Weyl semimetals have the two Weyl points in the Brillouin zone at \( \pm K_0 \) with \( K_0 = (0,0,k_0) \). Therefore, it is possible to consider two different configurations of the junction: the current parallel to \( z \) axis and the current parallel to \( x \) one. The characteristic features of the conductance spectra for the current parallel to \( z \) are essentially the same as those in the usual normal-metal/superconductor junctions. Namely, the conductance spectra becomes the bulk density of states in the superconductor when the normal transmission probability of the junction is low. In addition, the conductance spectra are insensitive to the weak magnetic moment at the junction interface. In this case, the chiral spin structure on the fermi surface does not affect the reflection process at the interface. In the case of the junction for the current parallel to the \( x \) axis, on the other hand, the chiral spin structure on the fermi surface suppress the Andreev reflection depending on the incident angles of a quasiparticle. This feature is explained by the spin mismatch between the incoming wave and the outgoing ones. The conductance spectra depends sensitively on the direction and the amplitudes of the magnetic moment at the interface because the spin flip scatterings relax the spin mismatch. The topological bound states appear on the \( xy \) surface of the Weyl semimetal. They, however, do not affect the low energy transport in the junctions.

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