Tunneling-conductivity anomalies at low bias voltages in a magnetic field

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The imposition of a magnetic field gives rise to structure at a value of eV equal in modulus to the Zeeman splitting for a conduction electron in addition to the anomalous behavior of the tunneling conductivity at a zero bias voltage V. The shape and height of these new structural features depend on the temperature, the spin scattering, and the details of the electron-electron interaction. The position of the structure is determined exclusively by the electron g-factor.

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Anomalies in the tunneling current at a zero bias voltage were first observed a long time ago and have now been observed in many experiments.[1] In many cases, the reasons for these anomalous features have remained puzzling for a long time.

The explanation offered in Ref. 2 for the tunneling anomalies has been shown by subsequent experiments[3–6] to successfully describe all the qualitative aspects of the phenomenon. Furthermore, the theory of Ref. 2 is in good quantitative agreement with experiment.

According to Ref. 2, a tunneling anomaly arises from a feature in the single-particle state density ν at the Fermi level at one of the electrodes (or simultaneously at both), caused by an interaction between conduction electrons. The energy dependence of the state density is determined by the effective dimensionality of the sample, d (Ref. 7):

$$\delta\nu_d = \frac{\lambda_\nu a_d}{(\hbar D)^{d/2}} \times \begin{cases} |\epsilon|^{\frac{d-2}{2}} & (d=1,3) \\ \ln|\epsilon|\tau/\hbar & (d=2) \end{cases}, \qquad (1)$$

$$a_1 = \frac{1}{4\sqrt{2\pi}}, \qquad a_2 = \frac{1}{4\pi^2}, \qquad a_3 = \frac{1}{8\pi^2\sqrt{2}} .$$

Here ε is the energy, reckoned from the Fermi level, D is the diffusion coefficient of the electrons, which is determined by their impurity scattering, τ is the mean free time $(\varepsilon \ll \hbar/\tau)$, ν_d is the d-dimensional state density, ν_2 is the number of states per unit area, and ν_1 is the number per unit length of the sample. The constant λ_ν is determined by the strength and nature of the electron-electron interaction.

The change in the tunneling conductivity G due to the increment (1) in ν_d is described at a sufficiently low temperature T by

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$$\frac{\delta G(V)}{G} = \frac{\delta \nu_d(\epsilon = eV)}{\nu_d} \qquad (eV \gg T), \tag{2}$$

where V is the bias voltage.

The effect of a magnetic field H on the state density and tunneling conductivity was discussed in part in Ref. 8, where it was shown that a magnetic field suppresses the effects of the interaction of electrons with a small resultant momentum (the "interaction effects in the Cooper channel"). With a short-range attraction of electrons, fluctuations of the superconducting order parameter give rise to an increment in the state density which has a minimum at $\varepsilon = 0$ if $T > T_c$, where T_c is the temperature of the transition to the superconducting state. The effect changes sign in the case of repulsion of the electrons. A magnetic field suppresses these features if $H > H_0 = c\varepsilon/4De$.

In this letter we examine the effect of a magnetic field on the correction to the state density for the interaction in the diffusion channel (which was studied in Ref. 2), and we show that in a magnetic field there are, in addition to the feature in the state density at $\varepsilon = 0$, some additional features (Zeeman features) at $\varepsilon = \pm\omega_S$, where $\omega_S = g\mu_B H$ (g is the gyromagnetic ratio of the conduction electron, and μ_B is the Bohr magneton). These features in the state density will be seen as anomalies in the tunneling conductivity at $eV = \pm\omega_S$. Since the condition $g\mu_B \ll 4De/c$ usually holds, Zeeman features will appear in fields at which the effects of the Cooper-channel interaction are completely suppressed.

In the diffusion channel, the two-particle Green's function is characterized by a momentum transfer $\mathbf{q}$ and an energy ω, which may be understood as the resultant momentum and resultant energy of the electron and hole. This Green's function is also characterized by the resultant spin of the electron and the hole, j, and its projections M. The interactions with different values of j and M contribute additively to $\delta\nu_d(\varepsilon)$:

$$\lambda_\nu = \lambda_\nu^{(j=0)} + 3\lambda_\nu^{(j=1)} .$$

In a magnetic field, the interaction with a given value of M gives rise to a feature in $\nu_d(\varepsilon)$ at $\varepsilon = -M\omega_S$. If $T = 0$, and there is no spin scattering, the correction to the state density is

$$\delta\nu_d(\epsilon) = \frac{a_d}{(\hbar D)^{d/2}} \left\{ [\lambda_\nu^{(j=0)} + \lambda_\nu^{(j=1)}] |\epsilon|^{\frac{d-2}{2}} + \lambda_\nu^{(j=1)} \left[|\epsilon + \omega_s|^{\frac{d-2}{2}} + |\epsilon - \omega_s|^{\frac{d-2}{2}} \right] \right\}$$

$$(d = 1, 3), \tag{3}$$

$$\delta\nu_2(\epsilon) = \frac{1}{4\pi^2 \hbar D} \left\{ [\lambda_\nu^{(j=0)} + \lambda_\nu^{(j=1)}] \ln\frac{|\epsilon|\tau}{\hbar} + \lambda_\nu^{(j=1)} \ln\frac{|\epsilon^2 - \omega_s^2|\tau^2}{\hbar^2} \right\} \quad (d=2) .$$

If (as is usually the case) $\lambda_\nu^{(j=0)} + \lambda_\nu^{(j=1)} > 0$, but $\lambda_\nu^{(j=1)} < 0$, then the imposition of a magnetic field will give rise to two maxima at $\varepsilon = \pm\omega_S$, in addition to the maximum at $\varepsilon = 0$, which becomes more intense than at $H = 0$.

The Zeeman structural features stem from the interaction of a particle from one spin subband, with an energy of approximately $2\alpha\omega_S$ ($\alpha = \pm 1/2$ is the projection of the spin in this subband onto the direction of **H**), with a particle near the Fermi level in the other subband, with nearly equal momenta (if the energy ε is not approximately equal to $2\alpha\omega_S$, the corresponding momentum is markedly different from the Fermi momentum in the other subband). In other words, the exact wave functions of the particles with $\varepsilon \approx 2\alpha\omega_S$ and $\varepsilon \cong 0$ from the different spin subbands have a strong spatial correlation.

The Zeeman structural features, in contrast with the basic feature at $\varepsilon = 0$, are spread out not only at a finite temperature but also by the spin scattering of conduction electrons. This scattering becomes important at

$$(\epsilon \pm \omega_s) \lesssim \hbar / t_s > T,$$

where t_S is the total scale time for spin relaxation, determined by both the spin-orbit scattering τ_{So} and the scattering by paramagnetic impurities,[9] τ_S:

$$t_s^{-1} = \frac{4}{3}\left(\tau_{so}^{-1} + \tau_s^{-1}\right).$$

The size and shape of the Zeeman structural features in the state density thus depend on the temperature, on the spin scattering, and—through the constant $\lambda_\nu^{(j=1)}$ —on the electron-electron interaction. The distance between these features along the energy scale, $2\omega_S$, however, is determined completely by the g-factor of the conduction electron.

It should also be noted that all the properties of the Zeeman structural features are independent of the ratio of ω_S and $\hbar/\tau$. Expression (3) remains valid near a feature with a given value of M even at $\omega_S > \hbar/\tau$, provided that $(\varepsilon + M\omega_S) \ll \hbar/\tau ((eV + M\omega_S) \ll \hbar/\tau)$.

The correction to the state density at $T = 0$ can be written

$$\delta\nu_d(\epsilon) = \sum_{j,M} \frac{\lambda_\nu^{(j)}}{2\pi} \mathrm{Re} \int (dq)\, D^{(j,M)}(\epsilon, \mathbf{q}), \tag{4}$$

where $(dq) = d^d q/(2\pi)^d$; an expression for the diffusion pole, $D^{(j,M)}$, with given values of j and M was derived in Ref. 10:

$$D^{(j,M)}(\omega, \mathbf{q}) = \left(-i\omega + Dq^2 - iM\omega_s \operatorname{sign}\omega + \frac{j}{t_s}\right)^{-1}. \tag{5}$$

Substitution of (5) into (4) yields an expression for the correction to the state density which becomes the same as (3) in the limit $t_S = \infty$. As mentioned earlier, the structural features resulting from the interaction with $j = 1$ spread out over a finite t_S. With $d = 3$, for example, this contribution can be found from (3) through the substitution

$$\sqrt{(\epsilon + M\omega_s)} \to \mathrm{Re}\sqrt{2\,[\,i(\epsilon + M\omega_s) + \hbar/t_s]} = \sqrt{\sqrt{(\epsilon + M\omega_s)^2 + \hbar^2/t_s^2} + \hbar/t_s},$$

where $M = 0, \pm 1$. For the interaction with $j = 0$, this contribution does not depend on t_S.

[1]J. M. Rowell, in: Tunnel'nye yavleniya v tverdykh telakh (Tunneling in Solids), Mir, Moscow, 1973, Ch. 27.
[2]B. L. Al'tshuler and A. G. Aronov, Solid State Commun. **30**, 115 (1979).
[3]W. L. McMillan and J. Mochel, Phys. Rev. Lett. **46**, 556 (1981).
[4]R. C. Dynes and J. P. Garno, Phys. Rev. Lett. **46**, 137 (1981).
[5]J. Imry and Z. Ovadyahu, Phys. Rev. Lett. **49**, 841 (1982).
[6]V. N. Bogomolov, E. V. Kolla, and Yu. A. Kumzerov, Materialy XXII Vsesoyuznogo soveshchaniya po fizike nizkikh temperatur (Proceedings of the XXII All-Union Conference on Low-Temperature Physics), Part II, Kishinev, 1982, p. 84.
[7]B. L. Al'tshuler, A. G. Aronov, and P. A. Lee, Phys. Rev. Lett. **44**, 1288 (1980).
[8]B. L. Al'tshuler, and A. G. Aronov, Solid State Commun. **38**, 11 (1981).
[9]B. L. Al'tshuler, A. G. Aronov, and A. Yu. Zuzin, Solid State Commun. **44**, 127 (1982).
[10]B. L. Al'tshuler, A. G. Aronov, and A. Yu. Zyuzin, Zh. Eksp. Teor. Fiz. **84**, No. 4 (1983) [Sov. Phys. JETP

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