Thermal Mechanism of Absolute Negative Conductivity in Two-Dimensional Electron Systems

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(Dated: March 22, 2022)

We calculate the dissipative dc conductivity of a two-dimensional electron system in a magnetic field for the situation when its effective temperature exceeds the temperature of the acoustic phonon system. We demonstrate that at sufficiently large difference between the electron and lattice temperatures (sufficiently strong microwave radiation) the dissipative dc conductivity can become negative. As an illustration, the case of the electron heating by microwave radiation is considered.

PACS numbers: PACS numbers: 73.40.-c, 78.67.-n, 73.43.-f

I. INTRODUCTION

Recent experimental studies of new types of magnetoresistance oscillations in a two-dimensional electron system (2DES) in a magnetic field [1, 2, 3, 4, 5], particularly the observation of vanishing electrical resistance caused by microwave radiation, have stimulated extensive efforts to clarify the origin of the uncovered phenomena. The occurrence of the effect of the so-called zero-resistance states is closely associated with the effect of absolute negative conductivity (ANC) when the dissipative dc conductivity \( \sigma \) is negative [6, 7, 8]. Mechanisms of ANC in a 2DES subjected to a magnetic field under nonequilibrium conditions, in particular under irradiation with microwaves, have been studied theoretically since the late 60s [9, 10, 11, 12, 13, 14, 15, 16, 17]. One can speculate that the photon-assisted impurity scattering can lead to the effect of ANC. The thermal mechanism of ANC causing the zero-resistance states. Notwithstanding this, a revision of alternate possible mechanisms of ANC may be of interest as well.

In this paper, we show that ANC can occur owing to the electron-phonon scattering (acoustical) if a 2DES is not in equilibrium with the phonon system. The deviation of the 2DES and the acoustic phonon systems from equilibrium can be caused, in particular, by the electron heating due to the absorption of microwaves or infrared radiation accompanied by the transitions between the Landau levels (LL’s) or the intersubband transitions. This is consistent with the concept considered previously [9, 10]. Here we derive the expression for the dc dissipative component of the 2DES conductivity \( \sigma \) as a function of the phonon (lattice) and 2DES (electron) temperatures, \( T \) and \( T_e \), respectively, considering the electron interaction with acoustic phonons. After that, we calculate \( T_e \) for the case of the electron heating by microwave radiation due to the absorption associated with the photon-assisted electron scattering on the phonons of the same type. We show that at sufficiently high electron temperatures (sufficiently strong radiation), the dissipative dc conductivity associated with the electron-phonon interaction can become negative, i.e., the electron heating can lead to the effect of ANC. The thermal mechanism of ANC under consideration can markedly influence the transport properties of a 2DES affected by microwave or infrared radiation.

II. CONDUCTIVITY

The dissipative dc current \( j_D(E) \) in a 2DES in the case of the electron-phonon interactions can be calculated from the following quite general expression:

\[
j_D(E) = \frac{e}{\hbar} \sum_{N, N'} f_N (1 - f_{N'}) \times \int d^3 q y_y |V_q|^2 |Q_{N, N'} (L^2 q_{\perp}^2 / 2)|^2 \\
\times \{ N_q \delta [(N - N') h\Omega_e + h\omega_q + eEL^2 y_y] \\
+ (N_q + 1) \delta [(N - N') h\Omega_e - h\omega_q + eEL^2 y_y] \}.
\] (1)

Here, \( L = (\hbar c/eH)^{1/2} \) is the quantum Larmor radius, \( \Omega_e = cH/m_e \) is the cyclotron frequency, \( f_N \) and \( N_q \), are the electron and phonon distribution functions, respectively, \( E \) is the net in-plane electric field, \( H \) is the magnetic field, \( N = 0, 1, 2, ... \) is the LL index, \( q = (q_x, q_y, q_z) \) (the directions \( x \) and \( y \) are in the 2DES plane), \( q_{\perp} = (q_x, q_y) \), and \( \omega_q = sq \) are the phonon wave vector, its in-plane component, and the phonon frequency, respectively, \( e = |e| \) is the electron charge, \( h \) is the Planck constant, \( c \) and \( s \) are the velocities of light and sound, respectively, \( \delta(\omega) \) is the form-factor of LL’s which at a small broadening \( \Gamma \) can be assumed to be the Dirac delta function, \( V_q \) is the matrix element of the electron-phonon interaction, \( |Q_{N, N'} (L^2 q_{\perp}^2 / 2)|^2 = |P_N^{N' N} (L^2 q_{\perp}^2 / 2)|^2 \exp(-L^2 q_{\perp}^2 / 2) \)

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is determined by the overlap of the electron wave functions, and \( |P_{N+1-N}^N(L^2q_{\perp}^2/2)|^2 \) is proportional to a Laguerre polynomial. Considering the acoustic piezoelectric scattering, we set \( |V_q|^2 \propto q^{-1} \exp(-l^2q_{\perp}^2/2), \) where \( l \) is the electron localization length in the \( z \)-direction perpendicular to the 2DES plane. Usually \( l \ll L \). We shall assume that the electron and phonon distributions are characterized by the respective temperatures \( T_e \) and \( T \): \( f_N = \exp(N\hbar\Omega_c - \zeta)/T_e^{-1} \), \( N' = \exp(\hbar\omega_q/T_1 - 1)^{-1} \), where \( \zeta \) is the Fermi energy reckoned from the lowest LL.

Assuming that \( eEL < h\Omega_c, \hbar\omega_q \), one can expand the expression in the right-hand side of Eq. (1) in powers of \( (eEL^2)/h\) and present the dissipative conductivity \( \sigma = f_D/E \) in the following form:

\[
\sigma = \left( \frac{e^2L^2}{\hbar^3c^2} \right) \sum_{N,A} \int d^2q q_{\perp}^2 |V_q|^2 \\
\times |Q_{N,N+1}(L^2q_{\perp}^2/2)|^2 \delta'(q - q^{(A)})
\]

\[
\times \left\{ f_N(1 - f_{N+1})N_q - f_{N+1}(1 - f_N)(N_q + 1) \right\}, \quad (2)
\]

where \( \delta'(q) = d\delta(q)/dq \), \( q^{(A)} = \Lambda\Omega_c/s, \) and \( \Lambda > 0 \). After replacing the integration over \( d^2q_{\perp} \) by the integration over \( dq_{\perp} dq_{\parallel} d\theta \), where \( \sin \theta = q_q/q_{\perp}, \) Eq. (3) can be reduced to

\[
\sigma \propto \sum_{N,A} \int_0^\infty dq_{\perp} q_{\perp}^3 \exp \left[ -\frac{l^2(q_{\perp}^2 - q_{\perp}^2)}{2} \right] \\
\times |Q_{N,N+1}(L^2q_{\perp}^2/2)|^2 \delta'(q - q^{(A)})
\]

\[
\times \left\{ f_N(1 - f_{N+1})N_q - f_{N+1}(1 - f_N)(N_q + 1) \right\}, \quad (3)
\]

Substituting the explicit formulas for the electron and phonon distributions functions into Eq. (3), one can thereafter present it as

\[
\sigma \propto \sum_{N,A} f_N(1 - f_{N+1}) \int_0^\infty dq \exp \left( -\frac{l^2q_{\perp}^2}{2} \right) G_N^{(A)}(q)
\]

\[
\times \left[ 1 - \exp(\hbar s q_{\perp}/T - \Lambda\Omega_c/T_e) \right] \delta'(q - q^{(A)}), \quad (4)
\]

where

\[
G_N^{(A)}(q) = \int_0^{L^2} dt t^3 \exp \left[ (t^2 - L^2)^2/2L^2 \right] |P_N^{(A)}(t^2/2)|^2.
\]

Assuming that \( T, T_e \ll \hbar\Omega_c \), and neglecting the terms containing higher powers of \( \exp(-h\Omega_c/kT) \), we arrive at

\[
\sigma \propto \int_0^\infty dq \exp(-l^2q_{\perp}^2/2) G_N(q)
\]

\[
\times \left[ \exp(-\hbar sq_{\perp}/T) - \exp(-\hbar\Omega_c/T_e) \right] \delta'(q - q^{(1)}), \quad (5)
\]

where \( G_N(q) = N_{N_m}(1 - f_{N_m+1})G_N^{(1)}(q) \) and \( N_m \) in the index of the LL immediately below the Fermi level. The integration in Eq. (5) gives

\[
\sigma \propto \exp(-h\Omega_c/T_e) \frac{d\exp(-l^2q_{\perp}^2/2) G_N(F(q))}{dq} \bigg|_{q=q^{(1)}}. \quad (6)
\]

Here

\[
F(q) = \frac{\exp(-h\Omega_c/T_e) - \exp(-h\Omega_c/T_e)}{\exp(-h\Omega_c/T_e)},
\]

At \( Lq \gg 1 \), one obtains \( G_N^{(A)}(q) \approx (1 - l^2/L^2)^{3/2} G_N^{(A)}/Lq \approx G_N^{(A)}/Lq \). If \( 1/\sqrt{N} < Lq < 1 \), one obtains \( G_N^{(A)}(q) \approx G_N^{(A)}(Lq)^2 \). Here \( G_N^{(A)} \) and \( G_N^{(A)} \) are coefficients. Taking into account that \( Lq^{(1)} = L\Omega_c/s \gg 1 \) and, hence, \( G_N^{(A)} \approx G_N^{(A)}/Lq \), at \( q \approx q^{(1)} \), from Eq. (6) we obtain

\[
\sigma = \sigma_0 = 1 + \left( \frac{T}{h\Omega_c} \right) \left( 1 + \frac{l^2\Omega_c^2}{s^2} \right)
\]

\[
\times \left[ 1 - \exp\left( \frac{h\Omega_c}{T} - \frac{h\Omega_c}{T_e} \right) \right]. \quad (7)
\]

Here \( \sigma_0 \) is the conductivity at \( T_e = T \) given by the following formula [12]:

\[
\sigma_0 \propto \left( \frac{h\Omega_c}{T L} \right)^2 \left( \frac{s}{L\Omega_c} \right) \exp\left( -\frac{h\Omega_c}{T} \right) \exp\left( -\frac{l^2\Omega_c^2}{2s^2} \right). \quad (8)
\]

At a small electron localization length \( l \), the last exponential factor in the right-hand side can be approximately replaced by unity.

Using Eq. (7), one can show that \( \sigma \) becomes negative when

\[
T_e > \frac{T}{1 - \theta} = T_1,
\]

where \( \theta = (T/h\Omega_c) \ln(h\Omega_c/T) \). The origin of ANC in such a situation is due to the following. When the electron temperature exceeds the lattice temperature, the balance between the processes of the phonon absorption accompanied by electron transitions to higher LL’s and the processes of the phonon emission accompanied by electron transitions to lower LL’s is violated, and the processes of the latter type dominate. The processes associated with the emission of a phonon with energy less
than the LL separation provide a negative contribution to the dissipative conductivity, whereas those with the energy exceeding this separation yield a positive contribution. Since the probability of the emission of a phonon with a lower energy is larger, the net contribution of the electron scattering processes involving the emission of phonons can be negative if the electron temperature is sufficiently large compared to the lattice temperature.

III. EFFECTIVE TEMPERATURE AT MICROWAVE HEATING

Let us consider the situation when the 2DES under consideration is heated by incoming microwave radiation. In this case, the electron distribution over LL’s is determined by the balance between the electron relaxation associated with the emission of acoustic phonons and the processes of the electron heating caused by the photon-assisted absorption/emission of such phonons. The balance equation can be presented in the following form:

\[ \frac{\hbar s}{N} \sum_{N} \sum_{\Lambda} f_N(1 - f_{N + \Lambda}) \int_0^\infty dq \, q \exp(-q^2/2)\tilde{G}_N^{(A)}(q) \times [\exp(-\hbar \Omega / T)(N_q + 1) - N_q] \delta(q - q^{(A)}) \]

\[ = \hbar \Omega J_{\Omega} \sum_{N} \sum_{\Lambda} f_N(1 - f_{N + \Lambda}) \int_0^\infty dq \, q \exp(-q^2/2)\tilde{G}_N^{(A)}(q) \times [N_q \delta(q + q^{(A)}) + (N_q + 1) \delta(q - q^{(A)})] \]

Here \( q^{(A)} = (\Omega - \Lambda \Omega_c)/s \), \( J_{\Omega} = (\tilde{E}_\Omega/\tilde{\Omega}_\Omega)^2 \), and \( \tilde{E}_\Omega \) and \( \tilde{\Omega}_\Omega \) are the microwave electric field amplitude and its characteristic value, respectively. The right-hand side of Eq. (10) is valid at sufficiently low microwave powers when the multi-photon processes (both real and virtual) are insignificant. At relatively low microwave powers when \( \tilde{E}_\Omega \ll \tilde{\Omega}_\Omega \), disregarding the polarization anisotropy, one can use the following formula (12) in (13):

\[ \tilde{\Omega}_\Omega = \tilde{\Omega}_\Omega \left( \Omega_c - \Omega \right) L/[e\Omega \sqrt{\Omega_c^2 + \Omega^2}] \]

where \( \tilde{\Omega}_\Omega = \sqrt{2m\Omega L/e} \).

At \( Lq \gg 1 \), the function \( K_N^{(A)}(q) \approx K_N^{(A)}/Lq \) with the coefficients \( K_N^{(A)} \) slightly different from \( G_N^{(A)} \). In the vicinity of the cyclotron resonance \( \Omega = \Omega_c \), the quantity \( \tilde{E}_\Omega \) is limited by the LL broadening if \( \tilde{E}_\Omega /\tilde{\Omega}_\Omega < \Gamma / \Omega_c \). In this limit one can put \( \tilde{E}_\Omega \approx \sqrt{2m\Omega \Gamma / e\tilde{\Omega}_\Omega \Gamma \Omega_c} \). In the case \( \tilde{E}_\Omega /\tilde{\Omega}_\Omega > \Gamma / \Omega_c \), when the multi-photon processes (both real and virtual) are essential, the situation becomes more complex (12).

Near the \( \Lambda \)th resonance \( |\Omega - \Lambda \Omega_c| \lesssim \Omega_c, T/h \ll \Omega_c \), Eq. (10) yields

\[ \frac{T_c - T}{T} \approx J_{\Omega} (\Omega - \Lambda \Omega_c)^2 \left( \frac{TL}{\hbar s^2} \right) \exp \left( \frac{\hbar \Omega}{T} \right) \Lambda \Phi_A. \]

Here

\[ \Phi_A = \tilde{M}_A^{(A)} \left( \frac{TL}{\hbar s^2} \right) \exp \left( \frac{L^2 \Omega_c^2}{2s^2} \right) \simeq \tilde{M}_A^{(A)} \]

and

\[ \tilde{M}_A^{(A)} = N_{n=0}^{N_{m}} \tilde{G}_N^{(A)} f_N(1 - f_{N + 1}) \left( \frac{N_{n=0}^{N_{m}}}{R_N^{(A)}} \right) f_N(1 - f_{N + 1}) \propto \Lambda. \]

One can see that the right-hand side of Eq. (11) tends to zero when the resonance detuning \( \Omega - \Lambda \Omega_c \) approaches zero. This occurs even at the cyclotron resonance \( \Lambda = 1 \) since an increase in \( J_{\Omega} \) at fixed microwave intensity is limited by the LL broadening: \( J_{\Omega} \propto [(\Omega - \Omega_c)^2 + \Gamma^2]^{-1} \). Therefore, near the cyclotron resonance

\[ \frac{T_c - T}{T} \propto \left( \frac{(\Omega - \Omega_c)^2}{(\Omega - \Omega_c)^2 + \Gamma^2} \right) \]

Thus, at the resonances, the electron heating associated with the photon-assisted acoustic scattering does not take place. In this case, the electron heating can occur due to other photon-assisted events. It is instructive that the dependence of the electron temperature on the resonance detuning \( \Omega - \Lambda \Omega_c \) is virtually an even function of the latter if \( |\Omega - \Lambda \Omega_c| \lesssim s/L, T/h \). When the resonance detuning is sufficiently large \( s/L, T/h \lesssim \Omega - \Lambda \Omega_c \lesssim \Omega_c \), the electron temperature becomes a markedly asymmetric function of \( \Omega - \Lambda \Omega_c \), particularly at very low lattice temperatures. Indeed, in such a range of detuning, from Eq. (10) we obtain the following relationships for the electron temperature:

\[ \exp \left( \frac{\hbar \Omega_c}{T} - \frac{\hbar \Omega_c}{T_c} \right) = 1 \]

\[ + \frac{J_{\Omega} \Omega}{\Omega - \Lambda \Omega_c} \exp \left( \frac{\hbar (\Omega - \Lambda \Omega_c)}{T} \right) \exp \left( \frac{\hbar \Omega_c}{T} \right) \Phi^{(A)} \]

at \( \Omega - \Lambda \Omega_c < 0 \), and

\[ \exp \left( \frac{\hbar \Omega_c}{T} - \frac{\hbar \Omega_c}{T_c} \right) = 1 \]

\[ + \frac{J_{\Omega} \Omega}{\Omega - \Lambda \Omega_c} \exp \left( \frac{\hbar \Omega_c}{T} \right) \Phi^{(A)} \]

at \( \Omega - \Lambda \Omega_c > 0 \), where \( \Phi^{(A)} \) is similar to \( \Phi^{(A)} \).

IV. SPECTRAL DEPENDENCIES

Using Eqs. (7) and (11) and assuming for simplicity that \( \hbar \Omega_c/s < 1, TL/\hbar s \approx 1, \tilde{\Phi}^{(A)} \simeq \Lambda, \Phi^{(A)} \simeq \Lambda, \)}
we obtain for the dissipative conductivity near the Λth resonance

\[
\sigma_0 - \sigma \simeq J_\Omega \frac{\Lambda^2 L^2 (\Omega - \Lambda \Omega_c)^2}{s^2} \exp \left( \frac{\hbar \Omega_c}{T} \right). \tag{15}\]

Using Eqs. (7) and (13), at the microwave frequencies well below the cyclotron resonance \(\Omega < \Omega_c\), we arrive at the following formula:

\[
\sigma_0 - \sigma \simeq \frac{J_\Omega}{|\Omega - \Omega_c|} \exp \left( \frac{\hbar (\Omega - \Omega_c)}{T} \right) \exp \left( \frac{-\hbar \Omega_c}{T} \right). \tag{16}\]

At higher microwave frequencies \(\Omega_c \gtrsim \Omega < 2\Omega_c\), Eqs. (7), (13), and (14) yield

\[
\frac{\sigma_0 - \sigma}{\sigma_0} \simeq \frac{J_\Omega}{|\Omega - \Omega_c|} \exp \left( \frac{\hbar \Omega_c}{T} \right). \tag{17}\]

The obtained dependences of \(\sigma\) versus \(\Omega/\Omega_c\) exhibit deep minima on the either side of the cyclotron resonances with \(\sigma < 0\) if the intensity of microwave radiation is sufficiently large. The positions of these minima are \(\Omega^\pm = \Omega_c \mp \delta_m\), with \(\delta_m \simeq s/L\). Similar minima, while less pronounced, occur at higher resonances. Comparing Eqs. (16) and (17), we obtain the following estimate:

\[|\sigma_0 - \sigma^+(\Omega)|/|\sigma_0 - \sigma^-(\Omega)| \simeq \exp(\hbar s/LT).\]

At sufficiently low lattice temperatures, this ratio is large because the contribution to the electron heating from the photon-assisted absorption of acoustic phonons is rather small. This implies that the dissipative dc conductivity at the microwave frequencies somewhat below the resonances (given by Eq. (7)) is close to its “dark” value \(\sigma_0\). The dependences of the normalized dissipative dc conductivity \(\sigma/\sigma_0\) on \(\Omega/\Omega_c\) calculated numerically using Eqs. (7) and (10) for different values of the parameter \(b = \hbar s/LT\) are shown in Fig. 1. For the numerical calculations we assumed that \(H = 2\) kG, \(l/L = 0.1\), \(\gamma = \Gamma/\Omega_c = 0.1\) and \(S = (E_0^2/2D)^2 = 10^{-3}\). Figure 2 shows the normalized dissipative conductivity \(\sigma/\sigma_0\) as a function of the parameter \(b \propto T^{-1}\) at different values of the detuning \(\Omega - \Omega_c\). One can see from Figs. 1 and 2 that at sufficiently large \(b\), i.e., at sufficiently low lattice temperatures, the ratio \(\sigma/\sigma_0\) as a function of \(\Omega\) and \(\Omega_c\) exhibits pronounced minima at which it can be negative.

Taking into account Eq. (17) and setting \(\Omega - \Omega_c = s/L\), one can obtain the following condition for ANC associated with the mechanism under consideration:

\[
J_\Omega > \left( \frac{s}{\Omega \Omega_c} \right) \exp \left( -\frac{\hbar \Omega_c}{T} \right). \tag{18}\]

The right-hand side of inequality (18) comprises two small parameters: \(s/L\Omega\) and \(\exp(-\hbar \Omega_c/T)\). For a 2DES in a GaAs/AlGaAs heterostructure, assuming \(s = 3 \times 10^5\) cm/s, \(H = 2\) kG, and \(T = 1\) K, the product of these parameters is estimated as \(10^{-3}\). Naturally, when the dark conductivity is determined by another mechanism than the acoustic phonon scattering, the occurrence of ANC requires a larger microwave power.

![FIG. 1: Normalized dissipative conductivity vs microwave frequency for different parameter \(b\) (inverse lattice temperature).](image1)

![FIG. 2: Normalized dissipative conductivity vs parameter \(b\) for different resonance detuning.](image2)

V. DISCUSSION

The electron interaction with interface acoustic phonons can also contribute to the mechanism of ANC under consideration. However, the dissipative conductivity associated with the electron scattering on 2D acoustic phonons is smaller than that associated with the scattering on 3D phonons by a large factor \(\exp((L^2 - l^2) \Omega_c^2/2s^2)\). Due to this, the variation of the dissipative conductivity caused by the electron heating is rather small. On the contrary, the contribution of the photon-assisted electron interaction with 2D phonons to the absorption of microwave radiation and, therefore, to the electron heating can be pronounced. The latter (extra) heating mechanism can result in an additional deepening of the conductivity minima.

Thus, the electron interaction with acoustic phonons can lead to ANC in a 2DES irradiated with microwaves. As shown in Refs. and here, two “acoustic” mechanisms of ANC are possible: (1) the dynamic mecha-
anism associated with the direct contribution of the photon-assisted scattering processes to the dissipative conductivity and (2) the thermal mechanism at which the microwave radiation increases the electron temperature that, in turn, affects the dissipative conductivity. Comparing the dynamic and thermal variations of the dissipative conductivity caused by irradiation, one can determine that their ratio does not include any parameter significantly different from unity. Hence, both the mechanisms in question can be simultaneously essential. However, these mechanisms lead to substantially different spectral behavior of the dissipative conductivity.

At the electron heating associated with the intersub-band transitions stimulated by infrared radiation, the dissipative conductivity can also become negative, exhibiting, however, a rather smooth (nonresonant) dependence on the magnetic field. The mechanism of ANC considered above is due to a difference between the the electron and phonon temperatures. Therefore, the dissipative conductivity in a 2DES subjected to a magnetic field can be negative not only at $T_e > T$ but at $T_e < T$ as well (see Ref. [10]).

VI. SUMMARY

We demonstrated a possibility of the thermal mechanism of negative dissipative dc conductivity in a 2DES subjected to a magnetic field associated with the interaction of electrons with acoustic phonons at a difference between the electron and phonon temperatures. As an example, the electron heating by microwave radiation is considered. It was shown that in this case, both dynamic and thermal mechanisms can be efficient to result in ANC at sufficiently strong microwave radiation in certain ranges of its frequency.

Acknowledgments

The author is grateful to V. A. Volkov and V. V. Vyurkov for useful discussions and A. Satou for assistance.

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