Fractional microwave-induced resistance oscillations

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Abstract
We develop a systematic theory of microwave-induced oscillations in magnetoresistivity of a 2D electron gas in the vicinity of fractional harmonics of the cyclotron resonance, observed in recent experiments. We show that in the limit of well-separated Landau levels the effect is dominated by the multiphoton inelastic mechanism. At moderate magnetic field, two single-photon mechanisms become important. One of them is due to resonant series of multiple single-photon transitions, while the other originates from microwave-induced sidebands in the density of states of disorder-broadened Landau levels.

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1. Introduction

Recently, much attention has been attracted to the discovery of microwave-induced resistance oscillations (MIRO) [1], followed by the spectacular observation of zero-resistance states (ZRS) in the oscillation minima [2,3]. Two microscopic mechanisms of the MIRO have been proposed: the “displacement” mechanism related to the effect of microwaves on the impurity scattering [4,5,6], and the “inelastic” mechanism accounting for nonequilibrium oscillatory changes in the electron distribution [7,8]. Both mechanisms rely on the energy oscillations of the density of states (DOS) ν(ε) of disorder-broadened Landau levels (LLs) and reproduce the observed phase of the ω/ωc-oscillations (here, ω and ωc = eB/mc are the microwave and the cyclotron frequencies). The inelastic mechanism yields temperature-dependent MIRO with the amplitude, proportional to the inelastic scattering time τi ∝ T−2, while the displacement contribution is T-independent, in disagreement with the experiments. At relevant T ∼ 1 K, the inelastic effect dominates and the corresponding theory [7,8] reproduces the experimental observations [1,2,3].

Further experimental investigations at elevated microwave power led to the discovery of “fractional” MIRO and ZRS [9,10] located near the fractional harmonics of the cyclotron resonance, ω/ωc = 1/2, 3/2, 5/2, 2/3, to be contrasted with the integer MIRO [1,2,3]. These remarkable observations motivated the present study where we address multiphoton effects and effects of the microwave radiation on the electronic spectrum, which govern the fractional MIRO in the case of separated LLs.

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2. Inelastic mechanism of the MIRO

We start by including the multiphoton processes in the theory [7,8]. In a classically strong magnetic field, \( \omega_c \tau_c \gg 1 \), the diagonal resistivity \( \rho_{xx} \) reads [7]

\[
\rho_{xx}/\rho_{xx}^D = \int d\varepsilon \bar{v}^2(\varepsilon) \partial_x f(\varepsilon),
\]

where \( \rho_{xx}^D = (e^2v_F^2\nu_0\tau_s)^{-1} \) is the Drude resistivity, \( v_F \) the Fermi velocity, \( \tau_s \) the transport scattering time, \( \nu_0 = m/2\pi \), and \( \bar{v}(\varepsilon) = \nu(\varepsilon)/\nu_0 \) the dimensionless density of states (DOS) of disorder-broadened LLs. The MIRO originate from the microwave-induced oscillations in the distribution function \( f(\varepsilon) \), which obeys the kinetic equation

\[
f(\varepsilon) - f_T(\varepsilon) = \frac{\tau_n}{4\tau_c} \sum_n A_n \bar{\nu}(\varepsilon-n\omega)[f(\varepsilon-n\omega) - f(\varepsilon)].
\]

Here \( A_n = A_{-n} \) describes the probability of \( n \)-photon absorption (emission) and \( f_T(\varepsilon) \) is the thermal distribution. The leading contribution to the integer MIRO comes from the single-photon \( A_1 = A_{-1} \equiv P_c \), where

\[
P_c = (\tau_n/\tau_c)(ev_F E_c^2)/\omega^2(\omega + \omega_c)^2.
\]

Here \( E_c \) is the amplitude of the circularly polarized microwave field [11], and \( \tau_n \ll \tau_c, \tau_s \) is the total (quantum) disorder-induced scattering time. If one assumes, in accord with the experimental conditions, that the temperature is high, \( 2\pi^2 T/\omega_c \gg 1 \), the contribution to \( \rho_{xx} \) of first order in \( P_c \) takes the form [7,8]

\[
\rho_{xx}/\rho_{xx}^D = \langle \bar{v}^2(\varepsilon) \rangle_c + (\tau_n/4\tau_c)P_c F(\omega),
\]

where \( \langle \ldots \rangle_c \) denotes the averaging over the period \( \omega_c \) of the DOS, and the function \( F(\Omega) \) oscillates with \( \Omega/\omega_c \),

\[
F(\Omega) = \Omega \langle \bar{v}^2(\varepsilon) \rangle_c \partial_c [\bar{\nu}(\varepsilon + \Omega) + \bar{\nu}(\varepsilon - \Omega)].
\]

In the limit of separated LLs, \( \omega_c \tau_c \gg 1 \), the DOS \( \bar{\nu}(\varepsilon) = \tau_{\lambda}\Ree[\sqrt{T^2 - (\varepsilon)^2}] \) is a sequence of semicircles of width \( 2T = 2(2\omega_c/\pi \tau_{\lambda})^{1/2} \), where \( \delta \varepsilon \) is the detuning from the center \( (n+\frac{1}{2})\omega_c \) of the nearest LL. In that case

\[
F(\Omega) = (16\Omega^2 \omega_c^2/3\pi^3 T^3) \Phi[(\Omega - N\Omega,\omega)/T],
\]

where \( N_0 \) is the integer number closest to \( \Omega/\omega_c \), and the odd function \( \Phi(x) \) is nonzero for \( |x| < 2 \) (Fig. 1)

\[
\Phi(x) = x(1 + |x|)\sqrt{|x|(2 - |x|) - 3x \arccos(|x| - 1)}.
\]

Provided \( \Gamma \ll \omega_c \), the oscillatory part of \( \rho_{xx} \) (4) is finite only in narrow intervals \( |\omega - (N + 1/2)\omega_c| < 4\Gamma \) around the integer values \( \omega/\omega_c = N \) [12]. Outside these intervals \( \nu(\varepsilon)\nu(\varepsilon + \omega) \equiv 0 \). Therefore, single-photon absorption is forbidden, so that \( f(\varepsilon) = f_T(\varepsilon) \) as long as multiphoton processes, given by \( A_n \) with \( |n| > 1 \), are not taken into account. Inclusion of two-photon processes leads to the appearance of the fractional MIRO in the frequency intervals \( |\omega - (N + 1/2)\omega_c| < 4\Gamma \). Outside these intervals \( \nu(\varepsilon)\nu(\varepsilon + 2\omega) \neq 0 \) (Fig. 2). In these intervals,

\[
\rho_{xx}/\rho_{xx}^D = \langle \bar{v}^2(\varepsilon) \rangle_c + (3\tau_n/32\tau_c)P_c^2 F(2\omega),
\]

where we used \( A_2 = A_{-2} = 3P_c^2/8 \) [6,13]. The doubling of the argument of the function \( F(\omega) \) in Eq. (7) [as compared to the integer MIRO, Eq. (4)] reflects the two-photon nature of the effect and leads to the emergence of the fractional MIRO at half-integer \( \omega/\omega_c \). The form and phase of the fractional oscillations (7) reproduce those for the integer MIRO, Eq. (4). Similarly to the integer case, there exists [14] a multiphoton contribution to the fractional MIRO governed by the displacement mechanism [4,5,6], which has a similar form but is a factor \( \omega_c \tau_n/\tau_c \Gamma \gg 1 \) smaller [16] than the inelastic one (7).
With increasing microwave power $P_\omega$, the resistivity (7) in the oscillation minima becomes negative, which indicates a transition to the ZRS [15]. Remarkably, like in the integer case [8], the leading-order approximation (7) for the multiphoton inelastic effect is sufficient to describe the fractional photoresponse even at such high power, since the second order contribution $\propto (P_\omega^2 / \tau_n / r_q)^2$ remains small in the parameter $\Gamma / \omega_c$.

4. Sideband mechanism of the fractional MIRO

Using the formalism developed in [6], it can be shown that the microwave illumination results in the appearance of "sidebands" in the DOS, located at distance $\omega$ on both sides of every LL (see Fig. 2). To first order in $P_\omega$ and assuming again $|\omega - (N + 1/2) \omega_c| < \Gamma$, we obtain the following expression for the microwave-induced sidebands [16]:

$$\tilde{\rho}^{(ab)}(\epsilon) = (\pi P_\omega / 8 \omega_c \tau_q) [\tilde{\rho}(\epsilon + \omega) + \tilde{\rho}(\epsilon - \omega)], \quad (8)$$

where $\tilde{\rho}(\epsilon)$ is the unperturbed DOS. In the presence of the sidebands, single-photon transitions become possible (Fig. 2), $\tilde{\rho}(\epsilon) \tilde{\rho}^{(ab)}(\epsilon \pm \omega) \neq 0$, resulting in the “sideband" contribution to the fractional MIRO,

$$\rho_{xx}^{(sb)} / \rho_{xx}^D = (\pi \tau_n / 64 \omega_c \tau_q^2) P_\omega^2 F(2\omega), \quad (9)$$

which has the same form as the leading two-photon inelastic contribution (7), but is smaller a factor $\omega_c \tau_q$. One more contribution originates from the sidebands oscillating in time with frequency $2\omega$. “Oscillating sideband" contribution [16] is symmetric with respect to the detuning from the fractional resonances, in contrast to the antisymmetric $F(2\omega)$, and is a factor $(\omega_c \tau_q)^{1/2}$ smaller than the two-photon contribution (7).

5. Conclusion

In the limit of well-separated LLs, $\omega_c \gg \Gamma$, the fractional MIRO are governed by the multiphoton inelastic mechanism, Eq. (7). At $\omega_c \sim \Gamma$, the sideband contribution (9) becomes relevant. Close to $\omega_c$, at which LLs start to overlap; specifically, at $\omega_c < 4\Gamma$, the effect is dominated by the resonant series of multiple single-photon transitions [10,17]. This effect appears at order $(\tau_n P_\omega / \tau_q)^2$. In the limit of strongly overlapping LLs, the fractional features get exponentially suppressed with respect to the integer MIRO [17,13].

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References

[1] M.A. Zudov, R.R. Du, J.A. Simmons, and J.R. Reno, Phys. Rev. B 64 (2001) 201311(R).
[2] R.G. Mani, J.H. Smet, K. von Klitzing, V. Narayanamurti, W.B. Johnson, and V. Umansky, Nature 420 (2002) 646.
[3] M.A. Zudov, R.R. Du, L.N. Pfeiffer, and K.W. West, Phys. Rev. Lett. 90 (2003) 046807.
[4] I.V. Gornyi, Sov. Phys. Solid State 11 (1970) 2078.
[5] A.C. Durst, S. Sachdev, N. Read, and S.M. Girvin, Phys. Rev. Lett. 91 (2003) 086803.
[6] M.G. Vavilov and I.L. Aleiner, Phys. Rev. B 69 (2004) 035303.
[7] I.A. Dmitriev, A.D. Mirlin, and D.G. Polyakov, Phys. Rev. Lett. 91 (2003) 226802.
[8] I.A. Dmitriev, M.G. Vavilov, L.L. Aleiner, A.D. Mirlin, and D.G. Polyakov, Phys. Rev. B 71 (2005) 115316.
[9] M.A. Zudov, R.R. Du, L.N. Pfeiffer, and K.W. West, Phys. Rev. B 73 (2006) 041303(R).
[10] S.I. Dorozhkin, J.H. Smet, K. von Klitzing, L.N. Pfeiffer, and K.W. West, cond-mat/0608633.
[11] Here we consider passive circular polarization of the microwave field. Qualitatively similar results for arbitrary polarization will be presented elsewhere [16].
[12] S.I. Dorozhkin, J.H. Smet, V. Umansky, and K. von Klitzing, Phys. Rev. B 71 (2005) 201306(R).
[13] I.A. Dmitriev, A.D. Mirlin, and D.G. Polyakov, Phys. Rev. B 75 (2007) 245320.
[14] X.L. Lei and S.Y. Liu, Appl. Phys. Lett. 88 (2006) 212109.
[15] A.V. Andreev, L.L. Aleiner, and A.J. Millis, Phys. Rev. Lett. 91 (2003) 056803.
[16] I.A. Dmitriev, A.D. Mirlin, and D.G. Polyakov, Phys. Rev. Lett. 99 (2007) 206805.
[17] I.V. Pechenezhskii, S.I. Dorozhkin, and I.A. Dmitriev, JETP Lett. 85 (2007) 86.