Temperature effects on microwave-induced resistivity oscillations and zero-resistance states in two-dimensional electron systems

J. Iñarrea[1,2] and G. Platero[1]
[1]*Instituto de Ciencia de Materiales, CSIC, Cantoblanco, Madrid 28049, Spain*
[2]*Escuela Politécnica Superior, Universidad Carlos III, Leganes, Madrid 28911, Spain*
(Received 8 July 2005; published 21 November 2005)

In this work we address theoretically a key issue concerning microwave-induced longitudinal resistivity oscillations and zero resistance states, namely, temperature. In order to explain the strong temperature dependence of the longitudinal resistivity and the thermally activated transport in two-dimensional electron gas, we have developed a microscopic model based on the damping suffered by the microwave-driven electronic orbit dynamics by interactions with the lattice ions yielding acoustic phonons. Recent experimental results show a reduction in the amplitude of the longitudinal resistivity oscillations and a breakdown of zero resistance states as the radiation intensity increases. In order to explain it we have included in our model the electron heating due to large microwave intensities and its effect on the longitudinal resistivity.

DOI: 10.1103/PhysRevB.72.193414 PACS number(s): 73.40.−c, 73.50.−h, 78.67.−n

Very few experiments in the field of condensed matter physics have produced such intense theoretical and experimental activities in the recent years as the one of longitudinal magnetoresisitvity (ρ_{xx}) oscillations and zero resistance states (ZRS).[1–4] These are obtained when a two-dimensional electron gas (2DEG) is subjected simultaneously to the influence of a moderate magnetic field (B) and microwave (MW) radiation. Many theoretical contributions have been presented to explain the physics behind and the dependence of ρ_{xx} with different variables such as MW intensity, frequency, and temperature (T).[5–10] Among all those contributions, very few of them have been devoted fully or partially to the study of the influence of T.[10–12] Experimental evidence[1–4] shows two common features concerning the dependence of ρ_{xx} with T: the first one is a reduction of ρ_{xx} oscillation amplitude as T is increased, eventually disappearing ZRS; the second one is the T variation of ρ_{xx} at the deepest minima which suggests thermally activated transport, that is, $\rho_{xx} \propto \exp(-E_{\rm act}/k_B T)$, where $E_{\rm act}$ is the activation energy.[1,2,4] In this paper we develop a microscopic model to explain most experimental results that involve the effect of T. In the first part we explain these features in terms of electron-phonon scattering. In the second part we include in our model the electron heating induced by the MW radiation. Then we can explain recently published experimental results[13,14] which demonstrate that at sufficiently high MW power, first the oscillatory amplitude becomes reduced by further increases in the MW power and second, a breakdown of ZRS is observed. We first study the influence of T through a damping parameter γ which affects dramatically the MW-driven electronic orbits harmonic movement:[10] along with this movement there occur interactions between electrons and lattice ions yielding acoustic phonons and producing a damping effect in the electronic motion. This is a *lattice* temperature (T_L) effect. We calculate γ through the electron-phonon scattering rate and the number of times that an electron interacts with lattice ions in its MW-driven harmonic motion. With this model we are able to explain not only the T_L dependence of ρ_{xx} but also why the T_L variation of the ρ_{xx} minima suggests a thermally activated transport. We can explain also the reduction of ρ_{xx} oscillation peak height along with an increase in the radiation intensity: we relate this effect with electron heating due to the corresponding increase of MW power. This is an *electron* temperature (T_e) effect.

Recently[10] we proposed a model to explain the ρ_{xx} behavior of a 2DEG at low B and under MW radiation. We obtained the exact solution of the corresponding electronic wave function

$$\Psi(x,t) = \phi_n[x - X - x_{\rm cl}(t),t] \times \exp\left[i\frac{m^*}{\hbar}\frac{dx_{\rm cl}(t)}{dt}[x - x_{\rm cl}(t)] + \frac{i}{\hbar}\int_0^t L dt' \right] \times \sum_{m=-\infty}^{\infty} J_m\left[\frac{eE_0}{\hbar} X\left(\frac{1}{w} + \frac{w}{\sqrt{(w_c^2 - w^2)^2 + \gamma^4}} \right) \right] e^{imwt}, \quad (1)$$

where e is the electron charge, ϕ_n is the solution for the Schrödinger equation of the unforced quantum harmonic oscillator, w is the MW frequency, w_c is the cyclotron frequency, E_0 is the intensity for the MW field, X is the center of the orbit for the electron motion, $x_{\rm cl}(t)$ is the classical solution of a forced harmonic oscillator, $x_{\rm cl} = [eE_0/m^*\sqrt{(w_c^2 - w^2)^2 + \gamma^4}]\cos wt$, L is the Lagrangian, and J_m are Bessel functions. According to that model, due to the MW radiation, center position of electronic orbits are not fixed, but they oscillate back and forth harmonically with w. The amplitude A for these harmonic oscillations is given by

$$A = \frac{eE_0}{m^*\sqrt{(w_c^2 - w^2)^2 + \gamma^4}}. \quad (2)$$

Now we introduce the scattering suffered by the electrons due to charged impurities randomly distributed in the sample. First, we calculate the electron-charged impurity scattering rate $1/\tau$ and second we find the average effective distance advanced by the electron in every scattering jump:

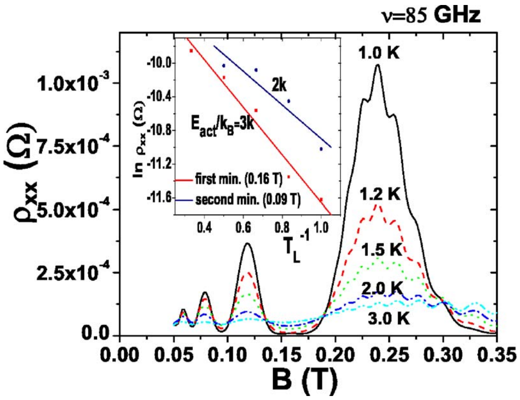

FIG. 1. (Color online) Calculated magnetoresistivity ρ_{xx} as a function of B, for different T_L at ν=85 GHz. In the inset we represent $\ln \rho_{xx}$ with T_L^{-1} at the first two deepest minima showing that $\rho_{xx} \propto \exp(-E_{act}/k_B T_L)$ reflecting thermally activated transport.

$\Delta X^{MW} = \Delta X^0 + A \cos w\tau$, where ΔX^0 is the effective distance advanced when there is no MW field present. Finally the longitudinal conductivity σ_{xx} can be calculated: $\sigma_{xx} \propto \int dE \Delta X^{MW}/\tau (f_i - f_f)$, where f_i and f_f are the corresponding distribution functions for the initial and final Landau states, respectively, and E is energy. To obtain ρ_{xx} we use the relation $\rho_{xx} = \sigma_{xx}/(\sigma_{xx}^2 + \sigma_{xy}^2) \simeq \sigma_{xx}/\sigma_{xy}^2$, where $\sigma_{xy} \simeq n_i e/B$ and $\sigma_{xx} \ll \sigma_{xy}$. At this point, we introduce a microscopic model which allows us to obtain γ and its dependence on T_L as follows. Following Ando and other authors,[15] we propose the next expression for the electron-acoustic phonons scattering rate valid at low T_L:

$$\frac{1}{\tau_{ac}} = \frac{m^* \Xi_{ac}^2 k_B T_L}{\hbar^3 \rho u_l^2 \langle z \rangle}, \tag{3}$$

where Ξ_{ac} is the acoustic deformation potential, ρ the mass density, u_l the sound velocity, and $\langle z \rangle$ is the effective layer thickness. However $1/\tau_{ac}$ is not yet the final expression for γ, this will be obtained multiplying $1/\tau_{ac}$ by the number of times that an electron in average can interact effectively with the lattice ions in a complete oscillation of its MW-driven back and forth orbit center motion. If we call n this number, then we reach the final expression $\gamma = 1/\tau_{ac} \times n$. An approximate value of n can be readily obtained in a simple way dividing the length an electron runs in a MW-induced oscillation (l_{osc}) by the lattice parameter of GaAs (a_{GaAs}): $n = l_{osc}/a_{GaAs}$. If w and w_c are approximately of the same order of magnitude, as it is in our case, l_{osc} turns out to be similar to the circular electronic orbit length. With the experimental parameters we have at hand[1] and for an average magnetic field it is straightforward to obtain a direct relation between γ and T_L, resulting in a linear dependence $\gamma\,(s^{-1}) \simeq 9.9 \times 10^{11}\,(s^{-1}\,K^{-1}) \times T_L$ (K). Now is possible to go further and calculate the variation of ρ_{xx} with T_L.

In Fig. 1, we present the calculated ρ_{xx} as a function of B for different T_L at $\nu = w/2\pi = 85$ GHz. In agreement with experiment,[1] as T_L is increased, the ρ_{xx} response decreases to eventually reach the darkness response. In the inset we represent the natural logarithm of ρ_{xx} with the inverse of T_L at the first two deepest minima showing that $\rho_{xx} \propto \exp(-E_{act}/k_B T_L)$, reflecting thermally activated transport. The qualitative behavior of ρ_{xx} as a function of B is very similar to the experimental one, however a quantitative agreement is still lacking. It could be due to the simplified model for the electronic scattering with impurities that we have considered.[10] In both experimental and calculated results, is surprising to see the different behavior of ρ_{xx} maxima and minima as T_L increases. In the first case, ρ_{xx} decreases and in the second one all the opposite, ρ_{xx} increases suggesting T_L-activated transport. We can find physical explanation as follows. In Fig. 2, we represent schematic diagrams to explain T_L dependence of maxima: due to the relation between T_L an γ, an increase in the first one means and increase in γ yielding a consequent reduction in the amplitude A [see formula (2)] of the MW-driven orbits oscillations. This has a dramatic impact in ΔX^{MW}, which is reduced as T_L is increased. A lesser

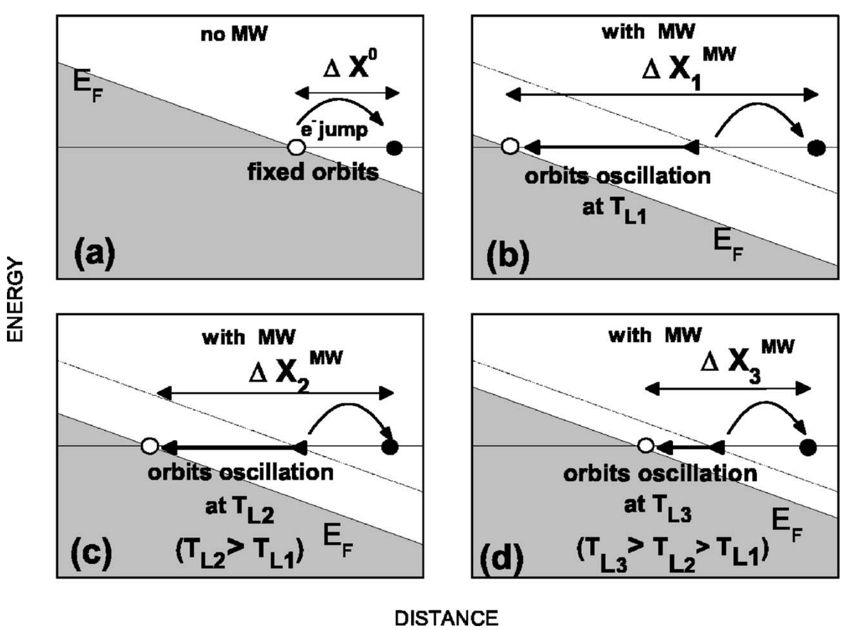

FIG. 2. Schematic diagrams of electronic transport without and with MW corresponding to *maxima* position. In (a) no MW field is present: due to scattering, electrons jump between fixed-position orbits and advance an effective average distance ΔX^0. When the MW field is on, the orbits are not fixed but oscillate with w. In (b)–(d) orbits move backwards during the jump, and on average electrons advance ΔX^{MW}, further than in the no MW case ΔX^0. As T_L becomes larger the orbits oscillation amplitude A is progressively reduced, reducing at the same time ΔX^{MW} and giving a progressive reduction in ρ_{xx} maxima.

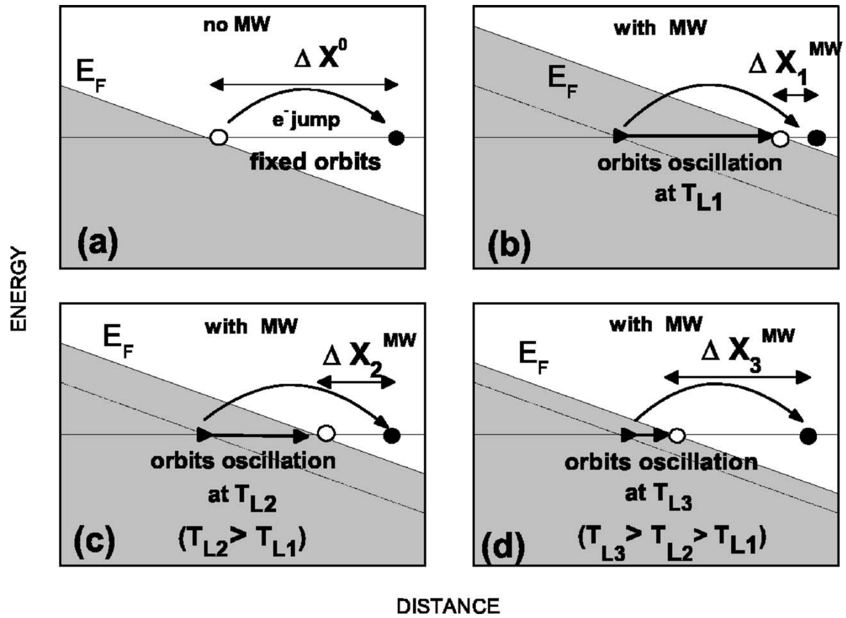

FIG. 3. Schematic diagrams of electronic transport corresponding to *minima* position. In this case, orbits move forward during the jump and on average electrons advance less than in the no MW case. Here, as in the case of maxima, an increase in T_L means a reduction in the orbits oscillation amplitude A, but for the case of minima this will produce a progressive larger ΔX^{MW} and consequently a larger conductivity and ρ_{xx}, which means a thermally activated transport. Then quenching of ZRS toward finite resistivity occurs. However if we increase the MW power, keeping constant T_L, A will be progressively larger. We can reach a situation where A is larger than the electronic jump, and the electron movement between orbits cannot take place because the final state is occupied. This situation corresponds to ZRS.

ΔX^{MW} [see Figs. 2(b)–2(d)] will imply a lesser conductivity and a progressive reduction of ρ_{xx} maxima. In Fig. 3 we can observe that, as in the case of maxima, an increase in T_L means a reduction in A, but for the ρ_{xx} minima this will produce a progressive larger ΔX^{MW} [see Figs. 3(b)–3(d)] and consequently a larger conductivity and ρ_{xx}, i.e., a thermally activated transport. Regarding E_{act}, now it is possible to explain its main functional dependencies as follows. We know that $E_{\mathrm{act}}/k_B = \Delta \ln \rho_{xx}/\Delta(1/T_L) \propto (\Delta\rho_{xx} \times \Delta T_L)$, where $\Delta\rho_{xx} = \rho_{xx}^{\mathrm{MW}}(T_L^{\mathrm{highest}}) - \rho_{xx}^{\mathrm{MW}}(T_L^{\mathrm{lowest}}) \simeq \rho_{xx}^{\mathrm{dark}} - \rho_{xx}^{\mathrm{MW}}(T_L^{\mathrm{lowest}})$. At increasing T_L the darkness response is eventually reached and $\rho_{xx}^{\mathrm{MW}}(T_L^{\mathrm{highest}}) \rightarrow \rho_{xx}^{\mathrm{dark}}$. In the case of a minimum with ZRS, $\rho_{xx}^{\mathrm{MW}}(T_L^{\mathrm{lowest}}) \rightarrow 0$. ΔT_L is the corresponding T_L difference. $\rho_{xx}^{\mathrm{dark}}$ is mainly sample dependent. It implies that $\Delta\rho_{xx}$ and therefore E_{act} will be also sample dependent and with different values for significantly different samples. If we consider the influence of MW power, it turns out that ρ_{xx}^{MW} is progressively smaller for increasing MW intensity yielding larger $\Delta\rho_{xx}$ and E_{act} as in experiments.[13] According to our model, the influence of ΔT_L on E_{act} is coming through γ and the consequent damping on the amplitude A. The mechanism of electron scattering responsible for the damping will be very important in the value of E_{act}. Thus, if the interaction is strong, the damping will be intense and A will be reduced very fast. Therefore ΔT_L and E_{act} will be small. However, if the interaction is weak, all the opposite will occur. We propose the electron-acoustic phonon interaction as the candidate to be responsible for the damping, reproducing most experimental features: linear behavior for $\ln \rho_{xx}$ vs $1/T_L$, value of ΔT_L and the order of magnitude of E_{act}.

Recent experiments[13,14] show that at sufficiently high MW intensities the ρ_{xx} oscillatory amplitude instead of getting larger, becomes reduced by further increases in the MW intensity. A breakdown of ZRS is also observed. We propose that it can be produced by electron heating occurring as the field intensity increases. To show that, we analyze theoretically the dependence of ρ_{xx} on T_e. It is clear that one of the effects of a progressive increase in the MW power will be electron heating and the corresponding increase in T_e. This will be reflected directly in the electronic distribution functions f_i and f_f, with the corresponding smoothing effect. Eventually we will obtain a progressive reduction in $(f_i - f_f)$ and therefore in ρ_{xx}. We find a situation where a MW power further increase yields two opposite effects on ρ_{xx} maxima. On the one hand a power rise will increase A [see

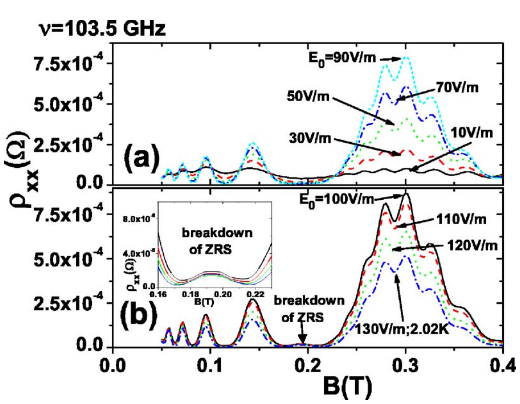

FIG. 4. (Color online) Magnetoresistivity ρ_{xx} response for different MW intensities. (a) Corresponds to lower intensities and (b) to higher ones. The inset shows amplified breakdown of ZRS. In both cases, $T_L = 1$ K.

formula (2)] which, in the maxima position, corresponds to larger ΔX^{MW} giving a ρ_{xx} rise. This is what is found in experimental and calculated results when the MW intensity is modest [see Fig. 4(a)]. On the other hand, a MW power rise yields also electron heating, increasing T_e ($E_0^2 \propto T_e^5$ according to available experimental evidence[16]) which, as we said above, will reduce the difference $(f_i - f_f)$ giving a reduction in ρ_{xx} maxima. This is what we find when the increasing MW power reaches a certain threshold value where the second effect is stronger than the first one, resulting in a progressive reduction of ρ_{xx} maxima [see Fig. 4(b)]. In Fig. 4(b) (see inset) we reproduce another surprising experimental result[13,14] as is the breakdown of ZRS states at high excitation power. Here at the minima, breakdown is characterized by a ρ_{xx} positive structure. Following our model, if we raise the MW intensity, we will eventually reach the situation where the orbits are moving forward but their amplitude is larger than the electronic jump, therefore $\Delta X^{\mathrm{MW}} < 0$. In that case the jump is blocked by Pauli exclusion principle $(f_i - f_f) = 0$, $\rho_{xx} = 0$, and we reach the ZRS regime. However at high MW intensities, in the ZRS regime, the MW-induced amplitude A is larger than the electronic jump and therefore the final state results to be below the Fermi energy. Under this regime, being $T_e > 0$ and due to the smoothing effect in the distribution function, we reach a situation where $f_f < 1$, $f_i < 1$, and $(f_i < f_f)$. Eventually we obtain that $(f_i - f_f)$ can be negative. In this situation $\Delta X^{\mathrm{MW}} < 0$ and $(f_i - f_f) < 0$ will produce an effective positive net current and positive ρ_{xx} in the middle of the ZRS region, resulting in the breakdown of this effect.

In summary, we have presented a theoretical model on the different effects of T on ρ_{xx} MW-driven oscillations and ZRS. The strong T_L dependence on ρ_{xx} and T_L-activated transport in 2DEG are explained through a microscopic model based on a damping process of the MW-driven orbits dynamics by interaction with acoustic phonons. Recent experimental results regarding a reduction in the amplitude of ρ_{xx} oscillations and breakdown of ZRS due to a further rise in the MW power are explained in terms of electron heating.

This work was supported by the MCYT (Spain) Grant No. MAT2002-02465, the "Ramon y Cajal" program (J.I.), and the EU Human Potential Programme HPRN-CT-2000-00144.

[1] R. G. Mani, J. H. Smet, K. von Klitzing, V. Narayanamurti, W. B. Johnson, and V. Umansky, Nature (London) **420**, 646 (2002).

[2] M. A. Zudov, R. R. Du, N. Pfeiffer, and K. W. West, Phys. Rev. Lett. **90**, 046807 (2003).

[3] S. A. Studenikin, M. Potemski, A. Sachrajda, M. Hilke, L. N. Pfeiffer, and K. W. West, cond-mat/0404411 (unpublished); S. A. Studenikin, M. Potemski, P. T. Coleridge, A. Sachrajda, and Z. R. Wasilewski, Solid State Commun. **129**, 341 (2004).

[4] R. L. Willett, L. N. Pfeiffer, and K. W. West, Phys. Rev. Lett. **93**, 026804 (2004).

[5] A. C. Durst, S. Sachdev, N. Read, and S. M. Girvin, Phys. Rev. Lett. **91**, 086803 (2003).

[6] X. L. Lei and S. Y. Liu, Phys. Rev. Lett. **91**, 226805 (2003).

[7] V. Ryzhii and V. Vyurkov, Phys. Rev. B **68**, 165406 (2003); V. Ryzhii, *ibid.* **68**, 193402 (2003).

[8] P. H. Rivera and P. A. Schulz, Phys. Rev. B **70**, 075314 (2004).

[9] A. V. Andreev, I. L. Aleiner, and A. J. Millis, Phys. Rev. Lett. **91**, 056803 (2003).

[10] J. Iñarrea and G. Platero, Phys. Rev. Lett. **94**, 016806 (2005); Microelectron. J. **36**, 334 (2005).

[11] S. A. Studenikin, M. Potemski, A. Sachrajda, M. Hilke, L. N. Pfeiffer, and K. W. West, cond-mat/0411338 (unpublished).

[12] X. L. Lei, J. Phys.: Condens. Matter **26**, 4045 (2004).

[13] R. G. Mani, Appl. Phys. Lett. **85**, 4962 (2004); Photonics Spectra **22**, 1 (2004); **25**, 189 (2004).

[14] S. A. Studenikin (private communication).

[15] T. Ando, A. Fowler, and F. Stern, Rev. Mod. Phys. **54**, 437 (1982); Karl Hess and C. T. Sah, Phys. Rev. B **10**, 3375 (1974); D. K. Ferry, *ibid.* **14**, 5364 (1976); C. T. Sah, T. H. Ning, and L. L. Tschopp, Surf. Sci. **32**, 561 (1972).

[16] R. R. Schliewe, A. Brensing, and W. Bauhofer, Semicond. Sci. Technol. **16**, 662 (2001).