Radiation-Induced Magnetoresistance Oscillation in a Two-Dimensional Electron Gas in Faraday Geometry

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Microwave-radiation induced giant magnetoresistance oscillations recently discovered in high-mobility two-dimensional electron systems are analyzed theoretically. Multiphoton-assisted impurity scatterings are shown to be the primary origin of the oscillation. Based on a theory which considers the interaction of electrons with electromagnetic fields and the effect of the cyclotron resonance in Faraday geometry, we are able not only to reproduce the correct period, phase, and the negative resistivity of the main oscillation, but also to predict the secondary peaks and additional maxima and minima observed in the experiments. These peak-valley structures are identified to relate, respectively, to single-, double-, and triple-photon processes.

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The discovery of a new type of giant magnetoresistance oscillations in a high-mobility two-dimensional (2D) electron gas (EG) subject to a crossed microwave (MW) radiation field and a magnetic field [1–5], especially the observation of "zero-resistance" states developed from the oscillation minima [4–7], has revived tremendous interest in magnetotransport in 2D electron systems [8–13]. These radiation-induced oscillations of longitudinal resistivity R_{xx} are accurately periodical in $1/B$, the inverse magnetic field, with the period determined by the MW frequency ω rather than the electron density N_e. The observed R_{xx} oscillations exhibit a smooth magnetic-field variation with resistivity maxima at $\omega/\omega_c = j - \delta_-$ and minima at $\omega/\omega_c = j + \delta_+$ (ω_c is the cyclotron frequency, $j = 1, 2, 3, \ldots$) having positive $\delta_\pm$ around $\frac{1}{4}$ [4]. The resistivity minimum goes downward with increasing sample mobility and/or increasing radiation intensity until a "zero-resistance" state shows up, while the Hall resistivity keeps the classical form $R_{xy} = B/N_e e$ with no sign of a quantum Hall plateau over the whole magnetic-field range exhibiting R_{xx} oscillation.

To explore the origin of the peculiar zero-resistance states, different mechanisms have been suggested [8–13]. It is understood that the appearance of negative longitudinal resistivity or conductivity in a uniform model suffices to explain the observed vanishing resistance [9]. The possibility of absolute negative photoconductance in a 2DEG subject to a perpendicular magnetic field was first explored 30 years ago by Ryzhii [14,15]. Recent works [8,10,11] indicated that the periodical structure of the density of states (DOS) of the 2DEG in a magnetic field and the photon-excited electron scatterings are the origin of the magnetoresistance oscillations. Durst *et al.* [8] presented a microscopic analysis for the conductivity assuming a δ-correlated disorder and a simple form of the 2D electron self-energy oscillating with the magnetic field, obtaining the correct period, phase, and the possible negative resistivity. Shi and Xie [11] reported a similar result using the Tien and Gorden current formula [16] for photon-assisted coherent tunneling. In these studies, however, the magnetic field is to provide an oscillatory DOS only and the high frequency (HF) field enters as if there is no magnetic field or with a magnetic field in Voigt configuration. The experimental setup requires us to deal with the magnetic field $\mathbf{B}$ perpendicular to the HF electric field. In this Faraday configuration, the electron moving due to HF field experiences a Lorentz force which gives rise to additional electron motion in the perpendicular direction. In the range of $\omega \sim \omega_c$, the electron velocities in both directions are of the same order of magnitude and are resonantly enhanced. This cyclotron resonance (CR) of the HF current response will certainly change the way the photons assist the electron scattering.

In this Letter, we construct a microscopic model for the interaction of electrons with electromagnetic fields in Faraday geometry. The basic idea is that, under the influence of a spatially uniform HF electric field, the center-of-mass (c.m.) of the whole 2DEG in a magnetic field performs a cyclotron motion modulated by the HF field of frequency ω. In an electron gas having impurity and/or phonon scatterings, there exist couplings between this c.m. motion and the relative motion of the 2D electrons. It is through these couplings that a spatially uniform HF electric field affects the relative motion of electrons by opening additional channels for electron transition between different states. Based on the theory for photon-assisted magnetotransport developed from this physical idea, we show that the main experimental results of the radiation-induced magnetoresistance oscillations can be well reproduced. We also obtain the secondary peaks and additional maxima and minima observed in the experiments [5,7].

For a general treatment, we consider N_e electrons in a unit area of a quasi-2D system in the x-y plane with a confining potential $V(z)$ in the z direction. These electrons, in addition to interacting with each other, are scattered by random impurities or disorders and by phonons in the lattice. To include possible elliptically polarized MW illumination, we assume that a uniform dc electric field $\mathbf{E}_0$ and ac field $\mathbf{E}_t \equiv \mathbf{E}_s \sin(\omega t) + \mathbf{E}_c \cos(\omega t)$ of frequency ω are applied in the x-y plane, together with a magnetic field $\mathbf{B} = (0, 0, B)$ along the z direction. In terms of the 2D c.m. momentum and coordinate of the electron system [17–19], which are defined as $\mathbf{P} \equiv \sum_j \mathbf{p}_{j\|}$ and $\mathbf{R} \equiv N_e^{-1} \sum_j \mathbf{r}_{j\|}$ with $\mathbf{p}_{j\|} \equiv (p_{jx}, p_{jy})$ and $\mathbf{r}_{j\|} \equiv (x_j, y_j)$ being the momentum and coordinate of the jth electron in the 2D plane, and the relative-electron momentum and coordinate $\mathbf{p}'_{j\|} \equiv \mathbf{p}_{j\|} - \mathbf{P}/N_e$ and $\mathbf{r}'_{j\|} \equiv \mathbf{r}_{j\|} - \mathbf{R}$, the Hamiltonian of the system can be written as the sum of a c.m. part $H_{\rm c.m.}$ and a relative-motion part $H_{\rm er}$ [$\mathbf{A}(\mathbf{r})$ is the vector potential of the $\mathbf{B}$ field]:

$$H_{\rm c.m.} = \frac{1}{2N_e m}[\mathbf{P} - N_e e\mathbf{A}(\mathbf{R})]^2 - N_e e(\mathbf{E}_0 + \mathbf{E}_t) \cdot \mathbf{R}, \quad (1)$$

$$H_{\rm er} = \sum_j \left[\frac{1}{2m}(\mathbf{p}'_{j\|} - e\mathbf{A}(\mathbf{r}'_{j\|}))^2 + \frac{p_{jz}^2}{2m_z} + V(z_j) \right] + \sum_{i<j} V_c(\mathbf{r}'_{i\|} - \mathbf{r}'_{j\|}, z_i, z_j), \quad (2)$$

together with the couplings of electrons to impurities and phonons, $H_{\rm ei}$ and $H_{\rm ep}$. Here m and m_z are, respectively, the electron effective mass parallel and perpendicular to the plane, and V_c stands for the electron-electron Coulomb interaction. Note that, although in $H_{\rm c.m.}$ and $H_{\rm er}$ the c.m. and the relative-electron motion are completely separated, the c.m. coordinate $\mathbf{R}$ enters $H_{\rm ei}$ and $H_{\rm ep}$ [18,19]. Starting from the Heisenberg operator equations for the rate of change of the c.m. velocity $\dot{\mathbf{V}} = -i[\mathbf{V}, H] + \partial \mathbf{V}/\partial t$ with $\mathbf{V} = -i[\mathbf{R}, H]$, and that of the relative-electron energy $\dot{H}_{\rm er} = -i[H_{\rm er}, H]$, we proceed with the determination of their statistical averages.

The c.m. coordinate $\mathbf{R}$ and velocity $\mathbf{V}$ in these equations can be treated classically, i.e., as the time-dependent expectation values of the c.m. coordinate and velocity [18], $\mathbf{R}(t)$ and $\mathbf{V}(t)$, such that $\mathbf{R}(t) - \mathbf{R}(t') = \int_{t'}^{t} \mathbf{V}(s)\, ds$. We are concerned with the steady transport under an irradiation of single frequency and focus on the photon-induced dc resistivity and the energy absorption of the HF field. These quantities are directly related to the time-averaged and/or base-frequency oscillating components of the c.m. velocity. At the same time, in an ordinary semiconductor the effect of higher harmonic current is safely negligible for the HF field intensity in the experiments. Hence, it suffices to assume that the c.m. velocity, i.e., the electron drift velocity, consists of a dc part $\mathbf{v}_0$ and a stationary time-dependent part

$$\mathbf{V}(t) = \mathbf{v}_0 + \mathbf{v}_1 \cos(\omega t) + \mathbf{v}_2 \sin(\omega t). \quad (3)$$

This time-dependent c.m. velocity enters all the operator equations having couplings to impurities and/or phonons in the form of the following exponential factor, which can be expanded in terms of Bessel functions $J_n(x)$:

$$e^{-i\mathbf{q}\cdot\int_{t'}^{t}\mathbf{V}(s)ds} = \sum_{n=-\infty}^{\infty} J_n^2(\xi) e^{i(\mathbf{q}\cdot\mathbf{v}_0 - n\omega)(t-t')} + \sum_{m\neq 0} e^{im(\omega t - \varphi)} \times \sum_{n=-\infty}^{\infty} J_n(\xi) J_{n-m}(\xi) e^{i(\mathbf{q}\cdot\mathbf{v}_0 - n\omega)(t-t')}.$$

Here $\xi \equiv \sqrt{(\mathbf{q}_\| \cdot \mathbf{v}_1)^2 + (\mathbf{q}_\| \cdot \mathbf{v}_2)^2}/\omega$ and $\tan\varphi = (\mathbf{q} \cdot \mathbf{v}_2)/(\mathbf{q} \cdot \mathbf{v}_1)$. On the other hand, for 2D systems having electron sheet density of the order of 10^{15} m^{-2}, the intraband and interband Coulomb interactions are sufficiently strong that it is adequate to describe the relative-electron transport state using a single electron temperature T_e. Except this, the electron-electron interaction is treated only in a mean-field level under random phase approximation (RPA) [18,19]. For the determination of unknown parameters $\mathbf{v}_0$, $\mathbf{v}_1$, $\mathbf{v}_2$, and T_e, it suffices to know the damping force up to the base-frequency oscillating term $\mathbf{F}(t) = \mathbf{F}_0 + \mathbf{F}_s \sin(\omega t) + \mathbf{F}_c \cos(\omega t)$, and the energy-related quantities up to the time-averaged term. We finally obtain the following force and energy balance equations:

$$0 = N_e e\mathbf{E}_0 + N_e e(\mathbf{v}_0 \times \mathbf{B}) + \mathbf{F}_0, \quad (4)$$

$$\mathbf{v}_1 = \frac{e\mathbf{E}_s}{m\omega} + \frac{\mathbf{F}_s}{N_e m\omega} - \frac{e}{m\omega}(\mathbf{v}_2 \times \mathbf{B}), \quad (5)$$

$$-\mathbf{v}_2 = \frac{e\mathbf{E}_c}{m\omega} + \frac{\mathbf{F}_c}{N_e m\omega} - \frac{e}{m\omega}(\mathbf{v}_1 \times \mathbf{B}), \quad (6)$$

$$N_e e\mathbf{E}_0 \cdot \mathbf{v}_0 + S_p - W = 0. \quad (7)$$

Here

$$\mathbf{F}_0 = \sum_{\mathbf{q}_\|} |U(\mathbf{q}_\|)|^2 \sum_{n=-\infty}^{\infty} \mathbf{q}_\| J_n^2(\xi) \Pi_2(\mathbf{q}_\|, \omega_0 - n\omega) + \sum_{\mathbf{q}} |M(\mathbf{q})|^2 \sum_{n=-\infty}^{\infty} \mathbf{q}_\| J_n^2(\xi) \Lambda_2(\mathbf{q}, \omega_0 + \Omega_{\mathbf{q}} - n\omega) \quad (8)$$

is the time-averaged damping force, S_p is the time-averaged rate of the electron energy gain from the HF field, $\frac{1}{2} N_e e(\mathbf{E}_s \cdot \mathbf{v}_2 + \mathbf{E}_c \cdot \mathbf{v}_1)$, which can be written in a form obtained from the right-hand side of Eq. (8) by replacing the $\mathbf{q}_\|$ factor with $n\omega$, and W is the time-averaged rate of the electron energy loss due to coupling with phonons, whose expression can be obtained from the second term on the right-hand side of Eq. (8) by replacing the $\mathbf{q}_\|$ factor with $\Omega_{\mathbf{q}}$, the energy of a wave vector $\mathbf{q}$ phonon. The oscillating frictional force amplitudes $\mathbf{F}_s \equiv \mathbf{F}_{22} - \mathbf{F}_{11}$ and $\mathbf{F}_c \equiv \mathbf{F}_{21} + \mathbf{F}_{12}$ are given by ($\mu = 1, 2$)

$$\mathbf{F}_{1\mu} = -\sum_{\mathbf{q}_\parallel} \mathbf{q}_\parallel \eta_\mu |U(\mathbf{q}_\parallel)|^2 \sum_{n=-\infty}^{\infty} [J_n^2(\xi)]' \Pi_1(\mathbf{q}_\parallel, \omega_0 - n\omega) - \sum_{\mathbf{q}} \mathbf{q}_\parallel \eta_\mu |M(\mathbf{q})|^2 \sum_{n=-\infty}^{\infty} [J_n^2(\xi)]' \Lambda_1(\mathbf{q}, \omega_0 + \Omega_{\mathbf{q}} - n\omega),$$

$$\mathbf{F}_{2\mu} = \sum_{\mathbf{q}_\parallel} \mathbf{q}_\parallel \frac{\eta_\mu}{\xi} |U(\mathbf{q}_\parallel)|^2 \sum_{n=-\infty}^{\infty} 2nJ_n^2(\xi) \Pi_2(\mathbf{q}_\parallel, \omega_0 - n\omega) + \sum_{\mathbf{q}} \mathbf{q}_\parallel \frac{\eta_\mu}{\xi} |M(\mathbf{q})|^2 \sum_{n=-\infty}^{\infty} 2nJ_n^2(\xi) \Lambda_2(\mathbf{q}, \omega_0 + \Omega_{\mathbf{q}} - n\omega).$$

In these expressions, $\eta_\mu \equiv \mathbf{q}_\parallel \cdot \mathbf{v}_\mu / \omega\xi$; $\omega_0 \equiv \mathbf{q}_\parallel \cdot \mathbf{v}_0$; $U(\mathbf{q}_\parallel)$ and $M(\mathbf{q})$ stand for effective impurity and phonon scattering potentials, $\Pi_2(\mathbf{q}_\parallel, \Omega)$ and $\Lambda_2(\mathbf{q}, \Omega) = 2\Pi_2(\mathbf{q}_\parallel, \Omega)[n(\Omega_{\mathbf{q}}/T) - n(\Omega/T_e)]$ [with $n(x) \equiv 1/(e^x - 1)$] are the imaginary parts of the electron density correlation function and electron-phonon correlation function in the presence of the magnetic field. $\Pi_1(\mathbf{q}_\parallel, \Omega)$ and $\Lambda_1(\mathbf{q}, \Omega)$ are the real parts of these two correlation functions. The effect of interparticle Coulomb interactions is included in them to the degree of level broadening and RPA screening.

The HF field enters through the argument ξ of the Bessel functions in $\mathbf{F}_0$, $\mathbf{F}_{\mu\nu}$, W, and S_p. Compared with that without the HF field ($n = 0$ term only) [20], we see that in an electron gas having impurity and/or phonon scattering (otherwise homogeneous), a HF field of frequency ω opens additional channels for electron transition: An electron in a state can absorb or emit one or several photons and is scattered to a different state with the help of impurities and/or phonons. The sum over $|n| \geq 1$ represents contributions of single and multiple photon processes of frequency-ω photons. These photon-assisted scatterings help to transfer energy from the HF field to the electron system (S_p) and give rise to additional damping force on the moving electrons. Note that $\mathbf{v}_1$ and $\mathbf{v}_2$ always exhibit CR in the range $\omega \sim \omega_c \equiv eB/m$, as can be seen from Eqs. (5) and (6) rewritten in the form

$$\mathbf{v}_1 = (1 - \omega_c^2/\omega^2)^{-1} \left\{ \frac{e}{m\omega} \left[\mathbf{E}_s + \frac{e}{m\omega} (\mathbf{E}_c \times \mathbf{B}) \right] + \frac{1}{N_e m\omega} \left[\mathbf{F}_s + \frac{e}{m\omega} (\mathbf{F}_c \times \mathbf{B}) \right] \right\}, \tag{9}$$

$$\mathbf{v}_2 = (\omega_c^2/\omega^2 - 1)^{-1} \left\{ \frac{e}{m\omega} \left[\mathbf{E}_c - \frac{e}{m\omega} (\mathbf{E}_s \times \mathbf{B}) \right] + \frac{1}{N_e m\omega} \left[\mathbf{F}_c - \frac{e}{m\omega} (\mathbf{F}_s \times \mathbf{B}) \right] \right\}. \tag{10}$$

Therefore, ξ may be significantly different from the argument of the corresponding Bessel functions in the case without a magnetic field or with a magnetic field in the Voigt configuration [20].

Equations (4)–(7) can be used to describe the transport and optical properties of magnetically biased quasi-2D semiconductors subject to a dc field and a HF field. Taking $\mathbf{v}_0 = (v_{0x}, 0, 0)$ in the x direction, Eq. (4) yields transverse resistivity $R_{xy} \equiv E_{0y}/N_e e v_{0x} = B/N_e e$, and longitudinal resistivity $R_{xx} \equiv E_{0x}/N_e e v_{0x} = -F_0/N_e^2 e^2 v_{0x}$. The linear magnetoresistivity is then

$$R_{xx} = -\sum_{\mathbf{q}_\parallel} q_x^2 \frac{|U(\mathbf{q}_\parallel)|^2}{N_e^2 e^2} \sum_{n=-\infty}^{\infty} J_n^2(\xi) \frac{\partial \Pi_2}{\partial \Omega} \Big|_{\Omega = n\omega} - \sum_{\mathbf{q}} q_x^2 \frac{|M(\mathbf{q})|^2}{N_e^2 e^2} \sum_{n=-\infty}^{\infty} J_n^2(\xi) \frac{\partial \Lambda_2}{\partial \Omega} \Big|_{\Omega = \Omega_{\mathbf{q}} + n\omega}. \tag{11}$$

The parameters $\mathbf{v}_1$, $\mathbf{v}_2$, and T_e in (11) should be determined by solving Eqs. (5)–(7) with a vanishing $\mathbf{v}_0$.

We calculate the unscreened $\Pi_2(\mathbf{q}_\parallel, \Omega)$ function of the 2D system in a magnetic field by means of Landau representation [17]:

$$\Pi_2(\mathbf{q}_\parallel, \Omega) = \frac{1}{2\pi l_B^2} \sum_{n,n'} C_{n,n'}(l_B^2 q_\parallel^2/2) \Pi_2(n, n', \Omega), \tag{12}$$

$$\Pi_2(n, n', \Omega) = -\frac{2}{\pi} \int d\varepsilon [f(\varepsilon) - f(\varepsilon + \Omega)] \times \mathrm{Im}\, G_n(\varepsilon + \Omega)\, \mathrm{Im}\, G_{n'}(\varepsilon), \tag{13}$$

where $l_B = \sqrt{1/|eB|}$ is the magnetic length, $C_{n,n+l}(Y) \equiv n![(n+l)!]^{-1} Y^l e^{-Y} [L_n^l(Y)]^2$ with $L_n^l(Y)$ the associate Laguerre polynomial, $f(\varepsilon) = \{\exp[(\varepsilon - \mu)/T_e] + 1\}^{-1}$ the Fermi distribution function, and $\mathrm{Im}\, G_n(\varepsilon)$ is the imaginary part of the Green's function, or the DOS, of the Landau level n. The real part functions $\Pi_1(\mathbf{q}_\parallel, \Omega)$ and $\Lambda_1(\mathbf{q}_\parallel, \Omega)$ can be obtained from their imaginary parts via the Kramers-Kronig relation.

We model the DOS function with a Gaussian-type form (ε_n is the energy of the nth Landau level) [21]:

$$\mathrm{Im}\, G_n(\varepsilon) = -(\pi/2)^{1/2} \Gamma^{-1} \exp[-(\varepsilon - \varepsilon_n)^2/(2\Gamma^2)], \tag{14}$$

with a broadening parameter $\Gamma = (2e\omega_c \alpha / \pi m \mu_0)^{1/2}$, where μ_0 is the linear mobility at temperature T in the absence of the magnetic field and $\alpha > 1$ is a semiempirical parameter to take account of the difference of the transport scattering time determining the mobility μ_0, from the single particle lifetime [4,8,10].

The moderate microwave intensity for the R_{xx} oscillation in these high-mobility samples yields only a slight electron heating, which is unimportant as far as the main phenomenon is concerned and is neglected for simplicity. We consider scatterings from remote impurities as well as from acoustic phonons. After solving $\mathbf{v}_1$ and $\mathbf{v}_2$ from Eqs. (9) and (10), the magnetoresisivity R_{xx} can be obtained directly from Eq. (11). At lattice temperature $T = 1$ K, the contribution from photon-assisted phonon

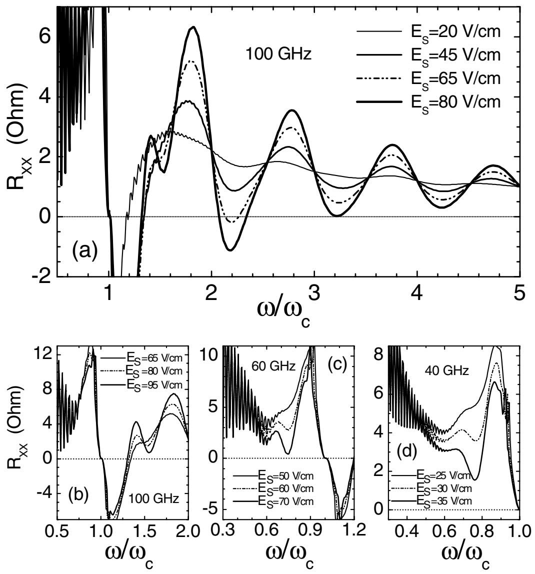

FIG. 1. The longitudinal magnetoresistivity R_{xx} of a GaAs-based 2DEG subject to a linearly polarized HF field $E_s \sin(\omega t)$. The parameters are temperature $T = 1$ K, electron density $N_e = 3.0 \times 10^{11}$ cm^{-2}, zero-magnetic-field linear mobility $\mu_0 = 2.4 \times 10^7$ cm^2 V^{-1} s^{-1}, and the coefficient $\alpha = 12$.

scattering is minor. The role of acoustic phonons, however, becomes essential at elevated lattice temperatures. Calculations were carried out for linearly polarized MW fields with multiphoton processes included.

Figure 1 shows the calculated longitudinal resistivity R_{xx} versus $\omega/\omega_c \equiv \gamma_c$ subject to a linearly polarized MW radiation of frequency $\omega/2\pi = 100$ GHz at four values of amplitude: $E_s = 20$, 45, 65, and 80 V/cm. Shubnikov–de Haas (SdH) oscillations show up strongly at the high ω_c side, and gradually decay away as $1/\omega_c$ increases. All resistivity curves exhibit pronounced oscillation having main oscillation period $\gamma_c = 1$ (they are crossing at points $\gamma_c =$ integer). The resistivity maxima locate around $\gamma_c = j - \delta_-$ and minima around $\gamma_c = j + \delta_+$ with $\delta_\pm \sim 0.23$–0.25 for $j = 3, 4, 5$, $\delta_\pm \sim 0.17$–0.21 for $j = 2$, and even smaller $\delta_\pm$ for $j = 1$. The magnitude of the oscillation increases with increasing HF field intensity for $\gamma_c > 1.5$. Resistivity gets into a negative value for $E_s = 80$ V/cm around the minima at $j = 1$, 2, and 3, for $E_s = 65$ V/cm at $j = 1$ and 2, and for $E_s = 20$ and 45 V/cm at $j = 1$. All these features, which are in fairly good agreement with experimental findings [1,3–5], are relevant mainly to single-photon ($|n| = 1$) processes. An anomaly appears in the vicinity of $\gamma_c = 1$, where the CR greatly enhances the effective amplitude of the HF field in photon-assisted scatterings and multiphoton processes show up. The amplitudes of the $j = 1$ maximum and minimum no longer monotonically change with field intensity. Furthermore, a shoulder appears around $\gamma_c = 1.5$ on the curves of $E_s = 45$ and 65 V/cm, and it develops into a secondary peak in the $E_s = 80$ V/cm case. This has already been seen in experiment (Fig. 2 of Ref. [5]). The valley between $\gamma_c = 1.4$ and 1.8 peaks can descend down to negative as E_s increasing further [Fig. 1(b)]. The appearance of the secondary peak is due to two-photon ($|n| = 2$) processes.

Radiation-induced resistivity behavior at $\gamma_c < 1$ is shown more clearly in the $\omega/2\pi = 60$ GHz case. As seen in Fig. 1(c), a shoulder around $\gamma_c = 0.4$–0.6 with a minimum at $\gamma_c = 0.6$ can be identified from the SdH oscillation background for all three curves, which is related to two-photon process. With increasing MW strength, there appear a clear peak around $\gamma_c = 0.68$ and a valley around $\gamma_c = 0.76$. This peak-valley is mainly due to three-photon ($|n| = 3$) process. In the case of 40 GHz, a similar peak and valley also show up [Fig. 1(d)].

We have also performed calculation using a Lorentz-type DOS function and find that, although the oscillating amplitude and the exact peak and valley positions are somewhat different, the basic feature of the radiation-induced magnetoresistivity oscillation remains.

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[1] M. A. Zudov *et al.*, Phys. Rev. B **64**, 201311(R) (2001).
[2] P. D. Ye *et al.*, Appl. Phys. Lett. **79**, 2193 (2001).
[3] R. G. Mani *et al.*, cond-mat/0305507.
[4] R. G. Mani *et al.*, Nature (London) **420**, 646 (2002).
[5] M. A. Zudov *et al.*, Phys. Rev. Lett. **90**, 046807 (2003).
[6] C. L. Yang *et al.*, Phys. Rev. Lett. **91**, 096803 (2003).
[7] S. I. Dorozhkin, JETP Lett. **77**, 681 (2003).
[8] A. C. Durst *et al.*, Phys. Rev. Lett., **91**, 086803 (2003).
[9] A. V. Andreev, I. L. Aleiner, and A. J. Millis, Phys. Rev. Lett. **91**, 056803 (2003).
[10] P. W. Anderson and W. F. Brinkman, cond-mat/0302129.
[11] J. Shi and X. C. Xie, Phys. Rev. Lett. **91**, 086801 (2003).
[12] J. C. Phillips, cond-mat/0303181.
[13] A. A. Koulakov, and M. E. Raikh, cond-mat/0302465.
[14] V. I. Ryzhii, Sov. Phys. Solid State **11**, 2078 (1970).
[15] V. I. Ryzhii, R. A. Suris, and B. S. Shchamkhalova, Sov. Phys. Semicond. **20**, 1299 (1986).
[16] P. K. Tien and J. P. Gordon, Phys. Rev. **129**, 647 (1963).
[17] C. S. Ting, S. C. Ying, and J. J. Quinn, Phys. Rev. B **14**, 5394 (1977).
[18] X. L. Lei and C. S. Ting, Phys. Rev. B **32**, 1112 (1985).
[19] X. L. Lei, J. L. Birman, and C. S. Ting, J. Appl. Phys. **58**, 2270 (1985).
[20] X. L. Lei, J. Appl. Phys. **84**, 1396 (1998); J. Phys. Condens. Matter **10**, 3201 (1998).
[21] T. Ando, A. B. Fowler, and F. Stern, Rev. Mod. Phys. **54**, 437 (1982).