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Interference oscillations of microwave photoresistance in double quantum wells

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We observe oscillatory magnetoresistance in double quantum wells under microwave irradiation. The results are explained in terms of the influence of subband coupling on the frequency dependent photoinduced part of the electron distribution function. As a consequence, the magnetoresistance demonstrates the interference of magnetointersubband oscillations and conventional microwave induced resistance oscillations.

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The phenomenon of microwave induced resistance oscillations[1–3] (MIRO) in two-dimensional (2D) electron systems under perpendicular magnetic fields has attracted much interest.[4] These oscillations are periodic in the inverse magnetic field with a period determined by the ratio of the microwave radiation frequency ω to the cyclotron frequency ω_c and survive at high temperatures. Basically, the observed oscillatory photoconductivity is caused by the Landau quantization of electron states, although different microscopic mechanisms of this phenomenon are still under discussion.

The influence of microwave irradiation on the magnetotransport properties is currently under investigation in quantum wells with a single occupied subband. In our Rapid Communication, we underline the importance of similar studies for the systems with two occupied 2D subbands, where the magnetotransport shows special features as compared to the single subband case. Apart from the commonly known Shubnikov-de Haas oscillations (SdHO), there exist the magnetointersubband (MIS) oscillations of resistivity caused by periodic modulation of the probability of intersubband transitions by the magnetic field[5,6] (see Ref. 7 for more references). These oscillations survive at high temperatures because they are not related to the position of the Landau levels with respect to the Fermi surface. Our recent observation[7] of large amplitude MIS oscillations at magnetic fields below 1 T in high mobility double quantum wells (DQWs) has established that these oscillations are a well reproducible feature of magnetotransport in such systems. Our present measurements, supported by a theoretical analysis, suggest that the behavior of oscillating magnetoresistance in DQWs under microwave irradiation is caused by an interference of the physical mechanisms responsible for the MIS oscillations and conventional MIRO.

We have studied dependence of the resistance of symmetric balanced GaAs DQWs with well widths of 14 nm and different barrier widths d_b=1.4, 2, and 3 nm on the magnetic field B in the presence of microwave irradiation of different frequencies and at different temperatures and intensities of radiation. The samples have high mobility of 10^6 cm^2/V s and high total sheet electron density $n_s \simeq 10^{12}$ cm^{-2}. The samples were mounted in a waveguide with different cross sections. The measurements were performed for perpendicular and parallel orientations of the current with respect to microwave polarization. The resistance $R = R_{xx}$ was measured by using the standard low-frequency lock-in technique. No polarization dependence of the resistance has been observed. While similar results were obtained for samples with different barrier widths, the data reported here correspond to 1.4 nm barrier samples. Without the irradiation, the magnetoresistance shows large period MIS oscillations clearly visible starting from B=0.1 T, and it also shows small period SdHO superimposed on the MIS oscillation pattern at higher fields[7] (see Fig. 1). The subband separation in our structure, Δ_{12}=3.67 meV, has been found from the MIS oscillation periodicity at low B. An increase in the intensity of the ra-

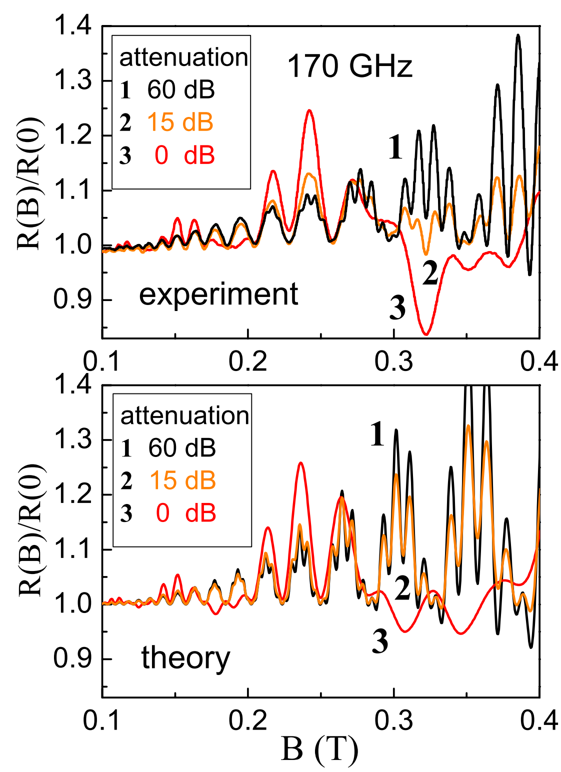

FIG. 1. (Color online) Measured (upper panel) and calculated (lower panel) magnetoresistance of DQWs for different intensities of 170 GHz radiation at T=1.4 K. The effect of the radiation can be neglected for the attenuation of 60 dB.

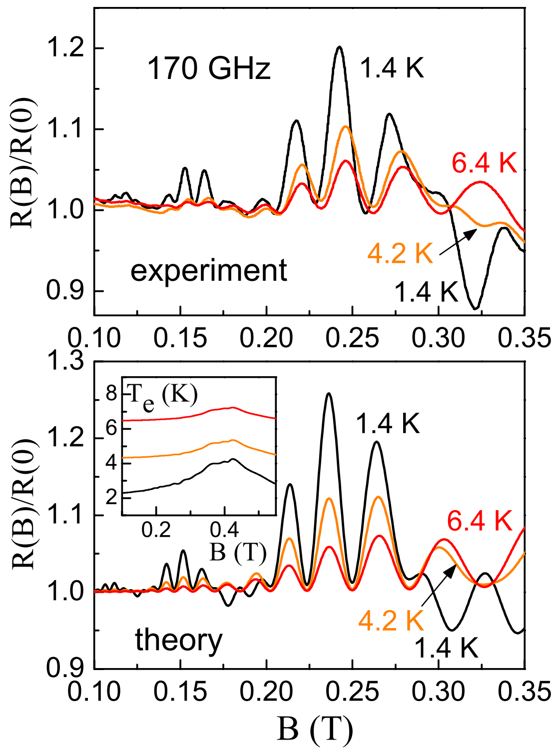

FIG. 2. (Color online) Measured (upper panel) and calculated (lower panel) magnetoresistance of DQWs under 170 GHz irradiation for different temperatures. The inset shows the dependence of electron temperature T_e on the magnetic field. Weak oscillations of T_e are caused by oscillations of the absorbed radiation power.

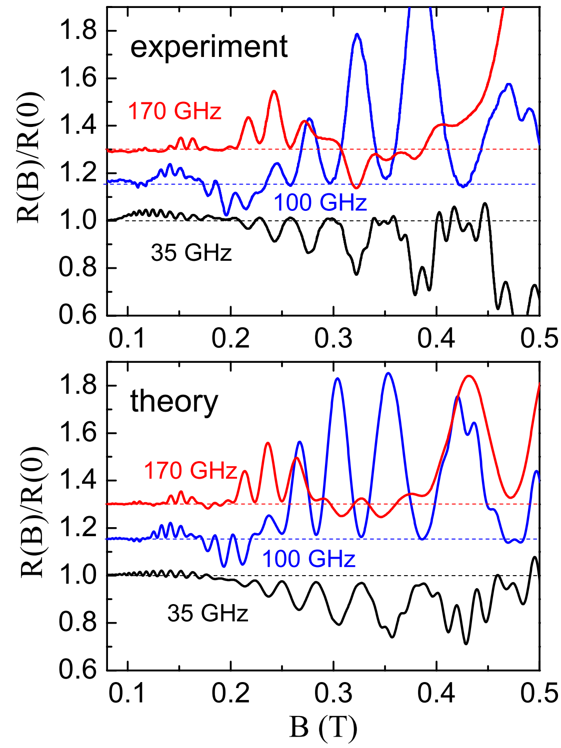

FIG. 3. (Color online) Measured (upper panel) and calculated (lower panel) magnetoresistance of DQWs at $T=1.4$ K for different frequencies of microwave excitation. The curves for 100 and 170 GHz are shifted up for clarity.

diation leads to the expected damping of the SdHO owing to electron heating and also causes suppression of some groups of MIS peaks, followed by the inversion (flip) of these peaks at high intensity, while the other MIS peaks are continuously enhanced by the irradiation. With increasing temperature, the inverted peaks go up, and the enhanced peaks go down, thus the conventional MIS oscillation picture is restored (Fig. 2). Measurements at different frequencies (Fig. 3) have confirmed that the selective peak flip is correlated with the radiation frequency and follows the periodicity determined by the ratio ω/ω_c. This allows us to attribute the observed effect to MIRO-related phenomena. However, the peak flip cannot be explained by a simple superposition of the factors $\cos(2\pi\Delta_{12}/\hbar\omega_c)$ and $-\sin(2\pi\omega/\omega_c)$ describing the MIS oscillations and the MIRO, respectively.

The theoretical explanation of our data is based on the physical model of Dmitriev *et al.*,[8] generalized to the two-subband case and improved by taking into account the electrodynamic effects in the absorption of radiation by 2D layers.[9–11] We assume that the elastic scattering of electrons, including the intersubband scattering, is much stronger than the inelastic one, so the electron distribution function (averaged over the time period $2\pi/\omega$) under microwave excitation remains quasi-isotropic and common for both subbands although essentially nonequilibrium. This energy distribution function, f_ε, is found from the kinetic equation $G_\varepsilon(f)=-J_\varepsilon(f)$, where G_ε is the generation term and J_ε is the inelastic collision integral. Under the excitation by linearly polarized high-frequency electric field $\mathbf{E}_t=\mathbf{E}\cos(\omega t)$, weak enough to neglect the multiphoton processes, the generation term is written as

$$G_\varepsilon = \frac{1}{D_\varepsilon}\frac{e^2}{8\pi m\omega^2}\sum_{\pm}|E_\omega^{(\pm)}|^2[\Phi_\varepsilon^{(\pm)}(\omega)(f_{\varepsilon+\hbar\omega}-f_\varepsilon) + \Phi_{\varepsilon-\hbar\omega}^{(\pm)}(\omega)(f_{\varepsilon-\hbar\omega}-f_\varepsilon)]. \tag{1}$$

where D_ε is the density of states, e is the electron charge, and m is the effective mass of electron. Next, $E_\omega^{(\pm)}=E\lambda_\omega^{(\pm)}$, where $\lambda_\omega^{(\pm)}=[1+2\pi\sigma_\pm(\omega)/c\sqrt{\epsilon}]^{-1}$, are the amplitudes of the circularly polarized components of electric field in the 2D plane, found from the Maxwell equations.[9,10] Here, $\sigma_\pm(\omega)$ are the complex conductivities describing the response of the electron system to the circularly polarized fields, c is the velocity of light, and ϵ is the dielectric permittivity of the medium surrounding the quantum well. The functions $\Phi_\varepsilon^{(\pm)}(\omega)$, which describe the probability of electron transitions between the states with energies ε and $\varepsilon+\hbar\omega$, are written in terms of Green's functions as $\Phi_\varepsilon^{(\pm)}(\omega)=\mathrm{Re}[Q_{\varepsilon,\omega}^{AR(\pm)}-Q_{\varepsilon,\omega}^{AA(\pm)}]$, where

$$Q_{\varepsilon,\omega}^{ss'(+)} = \frac{2\omega_c}{L^2}\sum_{nn'}\sum_{jj'}\sqrt{(n+1)(n'+1)}\sum_{p_y p_y'}\langle\langle G_\varepsilon^{jj',s}(n+1p_y,n'+1p_y')G_{\varepsilon+\hbar\omega}^{j'j,s'}(n'p_y',np_y)\rangle\rangle, \tag{2}$$

and $Q_{\varepsilon,\omega}^{ss'(-)}$ is given by the permutation of the indices, $n+1\leftrightarrow n$ and $n'+1\leftrightarrow n'$, in the arguments of Green's functions in Eq. (2). The Green's functions G, retarded (R) and advanced (A), are determined by the interaction of electrons with static disorder potential in the presence of magnetic field. They are written in the representation given by the product of the ket vectors $|j\rangle$ and $|np_y\rangle$, describing the 2D subbands and the Landau eigenstates, respectively. Here $j=1,2$ numbers the electron subband of the quantum well, n is the Landau level number, and p_y is the continuous momentum (the Landau gauge is used). The double angular brackets in Eq. (2) denote random potential averaging. The conductivities $\sigma_\pm(\omega)$ are expressed in terms of the functions (2) as

$$\sigma_\pm(\omega) = i\frac{e^2 n_s}{m\omega} + \frac{e^2}{2\pi m\omega}\int d\varepsilon[(f_\varepsilon - f_{\varepsilon+\hbar\omega})Q^{AR(\pm)}_{\varepsilon,\omega} + f_{\varepsilon+\hbar\omega}Q^{AA(\pm)}_{\varepsilon,\omega} - f_\varepsilon Q^{RR(\pm)}_{\varepsilon,\omega}]. \quad (3)$$

In the limit $\omega\to 0$, Eq. (3) gives the dc conductivity components $\sigma_d=\mathrm{Re}(\sigma_+ + \sigma_-)/2$ and $\sigma_\perp=\mathrm{Im}(\sigma_+ - \sigma_-)/2$.

Analytical evaluation of the correlation functions in Eq. (2) can be done in the case of relatively weak magnetic fields by using the self-consistent Born approximation and by expanding the Green's functions in powers of small Dingle factors. In application to balanced DQWs, where the subband energy separation is typically small compared to the Fermi energy, we assume that the difference in subband populations is small compared to n_s. It is justified for our samples where the Fermi energy (17 meV) is much larger than $\Delta_{12}/2=1.8$ meV. The difference between quantum lifetimes of electrons in the two subbands, τ_1 and τ_2, (and between the corresponding transport times) is also small, so we neglect it everywhere except the Dingle exponents. In the first order in the Dingle factors $d_j=\exp(-\pi/\omega_c\tau_j)$, the generation term acquires the following form:

$$G_\varepsilon = \mathcal{G}_\omega[r^{(0)}_\omega + (r^{(0)}_\omega - r^{(1)}_\omega)g_\varepsilon - r^{(1)}_\omega g_{\varepsilon+\hbar\omega}](f_{\varepsilon+\hbar\omega} - f_\varepsilon) + \{\omega \to -\omega\}, \quad (4)$$

where

$$\mathcal{G}_\omega = \frac{\pi e^2 E^2 n_s \tau_{tr}}{8m^2\omega^2}, \quad g_\varepsilon = \sum_{j=1,2} d_j \cos\frac{2\pi(\varepsilon-\varepsilon_j)}{\hbar\omega_c}, \quad (5)$$

$$r^{(0)}_\omega = \sum_\pm \frac{|\lambda^{(\pm)}_\omega|^2}{1+s_\pm}, \quad r^{(1)}_\omega = \sum_\pm \frac{|\lambda^{(\pm)}_\omega|^2 s_\pm}{(1+s_\pm)^2}, \quad (6)$$

ε_j are the subband energies, $s_\pm=(\omega\pm\omega_c)^2\tau^2_{\mathrm{tr}}$, and τ_{tr} is the transport time.

The kinetic equation is solved by representing the distribution function as a sum $f^0_\varepsilon + \delta f_\varepsilon$, where the first term slowly varies on the scale of cyclotron energy while the second one rapidly oscillates with energy.[8] The first term satisfies the equation $\mathcal{G}_\omega r^{(0)}_\omega(f^0_{\varepsilon+\hbar\omega}+f^0_{\varepsilon-\hbar\omega}-2f^0_\varepsilon)=-J_\varepsilon(f^0)$. Assuming that the electron-electron scattering controls the electron distribution, one can approximate f^0_ε by a heated Fermi distribution, $f^0_\varepsilon=\{1+\exp[(\varepsilon-\varepsilon_F)/T_e]\}^{-1}$. The electron temperature T_e is found from the balance equation $P_\omega=P_{\mathrm{ph}}$, where $P_\omega=\int d\varepsilon\varepsilon D_\varepsilon G_\varepsilon(f^0)=\mathrm{Re}[|\lambda^{(+)}_\omega|^2\sigma_+(\omega)+|\lambda^{(-)}_\omega|^2\sigma_-(\omega)]E^2/4$ is the power absorbed by the electron system and $P_{\mathrm{ph}}=-\int d\varepsilon\varepsilon D_\varepsilon J_\varepsilon(f^0)$ is the power lost to phonons. The second term satisfies the equation,

$$r^{(0)}_\omega[\delta f_{\varepsilon+\hbar\omega} + \delta f_{\varepsilon-\hbar\omega} - 2\delta f_\varepsilon] - \frac{\delta f_\varepsilon}{\tau_{in}\mathcal{G}_\omega} = r^{(1)}_\omega[g_{\varepsilon+\hbar\omega}(f^0_{\varepsilon+\hbar\omega} - f^0_\varepsilon) + g_{\varepsilon-\hbar\omega}(f^0_{\varepsilon-\hbar\omega} - f^0_\varepsilon)], \quad (7)$$

where we have used the relaxation-time approximation for the collision integral, $J_\varepsilon(\delta f)\simeq -\delta f_\varepsilon/\tau_{\mathrm{in}}$, with inelastic relaxation time τ_{in}. Solution of this equation is

$$\delta f_\varepsilon \simeq \frac{\hbar\omega_c}{2\pi}\frac{\partial f^0_\varepsilon}{\partial\varepsilon}\frac{A_\omega}{2}\sum_{j=1,2} d_j \sin\frac{2\pi(\varepsilon-\varepsilon_j)}{\hbar\omega_c}, \quad (8)$$

$$A_\omega = \frac{\mathcal{P}_\omega(2\pi\omega/\omega_c)\sin(2\pi\omega/\omega_c)}{1+\mathcal{P}_\omega\sin^2(\pi\omega/\omega_c)}\frac{r^{(1)}_\omega}{r^{(0)}_\omega}, \quad (9)$$

where $\mathcal{P}_\omega=4\mathcal{G}_\omega r^{(0)}_\omega\tau_{\mathrm{in}}$.

Using the energy distribution function found above, one can calculate the dc resistivity $\rho_{xx}=\sigma_d/(\sigma^2_d+\sigma^2_\perp)$:

$$\frac{\rho_d}{\rho_0} = 1 - 2\mathcal{T}_e g_{\varepsilon_F} + \frac{\tau_{\mathrm{tr}}}{\tau^{\mathrm{tr}}_{11}}(d_1^2+d_2^2) + 2\frac{\tau_{\mathrm{tr}}}{\tau^{\mathrm{tr}}_{12}}d_1 d_2\cos\frac{2\pi\Delta_{12}}{\hbar\omega_c} - \frac{A_\omega}{2}\left(d_1^2+d_2^2+2d_1d_2\cos\frac{2\pi\Delta_{12}}{\hbar\omega_c}\right), \quad (10)$$

where $\rho_0=m/e^2 n_s\tau_{\mathrm{tr}}$ is the zero-field Drude resistivity and $\Delta_{12}=\varepsilon_2-\varepsilon_1$ is the subband separation. The second term, proportional to g_{ε_F}, describes the SdHO. The third and the fourth terms describe positive magnetoresistance and the MIS oscillations,[7] respectively; these terms are written under the condition of classically strong magnetic fields, $\omega_c\tau_{\mathrm{tr}}\gg 1$. Here, $\tau^{\mathrm{tr}}_{11}=\tau^{\mathrm{tr}}_{22}$ and τ^{tr}_{12} are the intrasubband and intersubband transport scattering times, which contribute to the total transport time according to $1/\tau_{\mathrm{tr}}=1/\tau^{\mathrm{tr}}_{11}+1/\tau^{\mathrm{tr}}_{12}$. The term proportional to the oscillating factor A_ω describes modification of the oscillatory resistivity under photoexcitation. Another effect of the excitation is the electron heating, which increases thermal suppression of the SdHO described by the factor $\mathcal{T}_e=(2\pi^2T_e/\hbar\omega_c)/\sinh(2\pi^2T_e/\hbar\omega_c)$.

The most essential feature of the resistivity given by Eq. (10) is the presence of the product of the oscillating factors $\cos(2\pi\Delta_{12}/\hbar\omega_c)$ and $\sin(2\pi\omega/\omega_c)$, which corresponds to an interference of the MIS oscillations with the MIRO. The interference takes place because the photoexcitation involves both subbands [see Eqs. (8) and (9)], so the scattering induced coupling between the subbands, which oscillates as $\cos(2\pi\Delta_{12}/\hbar\omega_c)$, modifies the amplitude of the oscillating photocurrent. Usually, the microwave quantum energy is smaller than the subband separation, and the magnetoresistance shows fast MIS oscillations modulated by a slow MIRO component $\propto -\sin(2\pi\omega/\omega_c)$.

In our calculations, we have used Eq. (10) with $d_1=d_2=\exp(-\pi/\omega_c\tau)$ and $\tau^{\mathrm{tr}}_{12}=\tau^{\mathrm{tr}}_{11}=2\tau_{\mathrm{tr}}$, which is a good approximation for balanced DQWs.[7] We have taken into account the heating of electrons by the field and dependence of the characteristic times τ, τ_{tr}, and τ_{in} on the electron temperature T_e. The latter was determined by using the collision integral for interaction of electrons with both deformation and piezoelectric potentials of acoustic phonons. The inelastic scattering time τ_{in} is assumed to scale with the temperature as T_e^{-2}, with

$\hbar/\tau_{in}=4$ mK at $T_e=1$ K, according to the theoretical estimates for electron-electron scattering[8] applied to our samples. The temperature dependence of the transport time τ_{tr} and quantum lifetime τ of electrons in our samples has been empirically determined by analyzing temperature dependence of zero-field resistance and MIS oscillation amplitudes.[7] The microwave electric field $E \simeq 2.5$ V/cm is estimated by fitting calculated amplitudes of magnetoresistance oscillations to experimental data; the corresponding electron heating is in agreement with the observed suppression of the SdHO amplitudes.

The results of the calculations (shown in the lower parts of Figs. 1–3) are in good agreement with experiment and capture all qualitative features of the resistance dependence on the magnetic field, temperature, radiation power, and frequency. Note that the heating of electrons by microwaves is not strong at $\omega_c \simeq \omega$ (see inset to Fig. 2) because of radiative broadening of the cyclotron resonance[10,11] in our samples. The only feature that is not described by the theory is the observed reduction in the MIS oscillation frequency[7] at $B \gtrsim 0.25$ T, which is possibly related to a decrease in the subband separation owing to enhanced Coulomb correlations and/or modification of screening in magnetic field.

Understanding the MIRO physics requires comparison of measured photoresistance with the results provided by existing theories. Among the physical mechanisms[12] responsible for the MIRO, the inelastic mechanism[8,12] (describing the oscillations as a result of microwave induced change in the isotropic part of electron distribution function) is predicted to dominate at moderate radiation power—owing to a large ratio of τ_{in}/τ, which is of the order 10^2 under typical experimental conditions including our experiment. The observed insensitivity of the MIRO to the polarization of incident radiation and T^{-2} scaling of the oscillatory photoresistance amplitude[11] are in favor of the inelastic mechanism. We have shown that application of the basic principles of the inelastic mechanism theory[8] to the systems with two-subband occupation confirms the reliability of theoretical estimates for the inelastic relaxation time and leads to a satisfactory explanation of the features reported in this Rapid Communication.

In conclusion, we have presented experimental and theoretical studies of oscillatory magnetoresistance in DQWs under microwave irradiation. The interference oscillations observed in these systems appears because the photoinduced part of the electron distribution, which oscillates as a function of microwave frequency, is modified owing to subband coupling and becomes also an oscillating function of the subband separation.

Note added in proof. Recently, we have become aware of a related experimental work on a similar system.[13]

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