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Effect of density on microwave-induced resistance oscillations in back-gated GaAs quantum wells

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We report on microwave-induced resistance oscillations (MIRO) in a tunable-density 30 nm-wide GaAs/AlGaAs quantum well. We find that the MIRO amplitude increases dramatically with carrier density. Our analysis shows that the anticipated increase in the effective microwave power and quantum lifetime with density are not sufficient to explain the observed growth of the amplitude. We further observe that the fundamental oscillation extrema move towards cyclotron resonance with increasing density, which also contradicts theoretical predictions. These unexpected findings reveal that the density dependence is not properly captured by existing theories, calling for further studies.

Microwave-induced resistance oscillations (MIRO) appear in a two-dimensional (2D) electron gas (2DEG) [1–3] (or a 2D hole gas [4, 5]) subjected to low temperature $T$, weak magnetic field $B$, and microwave radiation of frequency $f = \omega/2\pi$. It is well established experimentally that MIRO originate from the bulk of the 2DEG [6–16]. Theoretically, microwave photoresistance $\delta R$ due to MIRO can be described by [17–19]

$$\frac{\delta R(\epsilon)}{R_0} = -2\pi \epsilon^2 \lambda^2 P \eta \sin 2\pi \epsilon.$$  (1)

Here, $R_0$ is the resistance at $B = 0$, $\epsilon = \omega/\omega_c$, $\omega_c = eB/m^*$ is the cyclotron frequency of the charge carrier with the effective mass $m^*$, $\lambda = \exp(-\epsilon/2f\tau_q)$ is the Dingle factor, $\tau_q$ is the quantum lifetime, $P$ is the effective microwave power, and $\eta$ is the dimensionless parameter (discussed later in detail) which depends on the disorder characteristics and the inelastic relaxation. The above expression was obtained assuming $2\pi k_BT \gg h\omega$ and is accurate away from the cyclotron resonance ($2\pi \epsilon \gg 1$), when the microwave power is not too high ($P \ll 1$), and when Landau levels are overlapping ($\lambda \ll 1$).

To date, MIRO have been observed in three kinds of material systems, namely GaAs/AlGaAs [1, 2], Ge/SiGe [4, 5], and MgZnO/ZnO [3] heterostructures, and the dependence on $\epsilon$, $\delta \rho \propto -e\lambda^2 \sin 2\pi \epsilon$, has been verified in many experiments. In addition, it was established that while MIRO at high microwave power significantly deviate from Eq. (1), they can still be well described within the same theoretical framework after generalization to arbitrary radiation intensity [10, 20, 21]. At the same time, experiments also revealed situations when existing theory is inadequate, e.g., in describing the measured dependencies on radiation polarization [22, 23] and on temperature [24]. Limitations of the theory were also identified in the regime of separated Landau levels [25] and in the radiation-induced modification of Shubnikov-de Haas oscillations [26].

One important parameter, role of which remained largely unexplored, is the carrier density $n_e$. While it has been recently demonstrated that $n_e$ affects $\epsilon$, presumably through interaction-induced renormalization of the effective mass $m^*$ [27, 28], it should also modify other parameters, e.g., $P$ and $\eta$, entering Eq. (1). Since the density dependencies of $P$ and $\eta$ are both known theoretically, MIRO measurements as a function of $n_e$ should provide an important test to existing microscopic description of microwave photoresistance.

In this Rapid Communication we investigate the effect of the carrier density on the MIRO amplitude employing a tunable-density 2DEG [28–31]. We find that the quantum lifetime depends on the carrier density only weakly, in agreement with the recent study investigating Shubnikov-de Haas oscillations in a similar device [31]. Our main finding, however, is a significant growth of the MIRO amplitude with the carrier density. The analysis shows that this growth cannot be accounted for by anticipated density dependence of $P\eta$ entering the prefactor of Eq. (1) [17–19, 32, 33]. Furthermore, we find that the fundamental extrema move towards $\epsilon = 1$ which also contradicts theoretical expectations. Both of these findings indicate that our understanding of the role of density in microwave photoresistance remains limited calling for further investigations.

Our 2DEG resides in a 30-nm GaAs/AlGaAs quantum well located about 200 nm below the sample surface. The structure is doped in a 2 nm GaAs quantum well at a setback of 63 nm on a top side. The in situ back gate consists of an $n^+$ GaAs layer situated 850 nm below the bottom of the quantum well. Ohmic contacts were fabricated at the corners and midsides of the lithographically-defined $1 \times 1 \text{mm}^2$ Van der Pauw mesa. The density of the 2DEG was varied from $n_e = 1.41 \text{ to } 2.87 \times 10^{11} \text{ cm}^{-2}$. Over this density range, the low-temperature elec-
electron mobility increased from \( \mu \approx 0.37 \) to \( \mu \approx 1.1 \times 10^7 \) cm\(^2\)V\(^{-1}\)s\(^{-1}\) [34], roughly following \( \mu \approx n_e^\alpha \), with \( \alpha \approx 1.5 \) [35, 36]. Microwave radiation, generated by a synthesized sweeper, was delivered to the sample immersed in liquid \(^3\)He via a rectangular (WR-28) stainless steel waveguide. The resistance \( R \) was measured using a standard low-frequency (a few Hz) lock-in technique.

In Fig. 1(a) we present \( R/R_0 \) vs \( 1/\epsilon \) for three different densities, \( n_e \approx 1.41 \) (bottom trace), 2.15 (middle trace), and 2.87 \times 10^{11} \) cm\(^{-2}\) (top trace), measured at \( T = 1.5 \) K under irradiation by microwaves of \( f = 34 \) GHz. The data show that MIRO become significantly stronger with increasing density. We next extract the oscillatory correction \( \delta R/R_0 \) and present the result in Fig. 1(b) as a function of \( \epsilon \). All three data sets reveal expected periodicity with \( \epsilon \), in agreement with Eq. (1).

According to Eq. (1), the MIRO amplitude is proportional to the effective microwave power, which was shown to be [32, 38]

\[
\mathcal{P}(\epsilon) = \frac{P_0}{2} \sum_{\pm} \frac{1}{(1 \pm \epsilon^{-1})^2 + \beta^2}, \quad P_0 = \frac{e^2 \varepsilon_{ac}^2 v_F^2}{\epsilon_{\text{eff}} h^2 \omega^4}. \tag{2}
\]

Here, \( \beta_\omega \equiv (\omega \tau_{\text{em}})^{-1} + (\omega \tau)^{-1}, \tau = (m^*/\epsilon)\mu \) is the momentum relaxation time, \( \tau_{\text{em}}^{-1} = n_e e^2/2\sqrt{\epsilon \epsilon_{\text{eff}} m^* c} \) [38] is the radiative decay rate, \( 2\sqrt{\epsilon \epsilon_{\text{eff}}} = \sqrt{\epsilon} + 1 \) defines the effective dielectric constant \( \epsilon_{\text{eff}} \), \( \epsilon = 12.8 \) is the dielectric constant of GaAs, \( v_F \) is the Fermi velocity, and \( \varepsilon_{ac} \) is the microwave electric field. The density dependence of \( \beta_\omega \) has been recently verified in time-resolved measurements of the cyclotron resonance [33]. Within the density range studied in our experiment, \( \tau \gg \tau_{\text{em}} \) and \( \beta_\omega \approx (\omega \tau_{\text{em}})^{-1} \propto n_e \). However, \( \beta_\omega \) remains much smaller than unity and, as result, \( \mathcal{P} \) increases with \( n_e \) for all \( \epsilon \) except in close proximity to \( \epsilon = 1 \). We will see, however, that the anticipated increase in \( \mathcal{P} \) is rather small and contributes little to the growth of MIRO shown in Fig. 1.

The growth of MIRO with \( n_e \) observed in Fig. 1 can, in principle, stem from \( \tau_q \) (entering \( \lambda \)) or \( \eta \). Both of these parameters are readily available from the Dingle analysis. Following Eq. (1), we introduce a reduced MIRO amplitude \( A \equiv |\delta R|_{\text{max}}/2\pi c P R_0 \) [39], where \( |\delta R|_{\text{max}} \) is the MIRO amplitude, and present it in Fig. 2 as a function of \( \epsilon \) for \( n_e \approx 1.41 \) (circles), 2.15 (squares), and 2.87 \times 10^{11} \) cm\(^{-2}\) (triangles). Fits to the data with \( A_0 \exp(-\epsilon/f \tau_q) \) (solid lines) yield \( \tau_q \approx 19.5, 20.0, \) and 21.2 ps respectively, indicating a slight increase of \( \tau_q \) with \( n_e \). In contrast, the intercept of the Dingle plots, \( A_0 \), grows substantially with \( n_e \). As we show below, theory predicts that under our experimental conditions \( A_0 \) can only decrease with \( n_e \).

After repeating the Dingle analysis for other \( n_e \), we present the density dependence of \( \tau_q \) (circles) in Fig. 3. A slight increase of \( \tau_q \) with \( n_e \) appears to contradict a recent study [31], which has found a saturation of \( \tau_q \) at \( n_e \approx 2 \times 10^{11} \) cm\(^{-2}\) and a monotonic decrease at higher \( n_e \). This discrepancy can be alleviated by recalling that Shubnikov-de Haas oscillations employed in Ref. 31 yield only impurity contribution to the quantum lifetime, \( \tau_{\text{q0}} \) [40, 41]. The quantum lifetime obtained from MIRO, on the other hand, is reduced by electron-electron scatter-
increasing in the range studied. This finding is unexpected since, as we noted from the data in Fig. 1(b), these extrema move to slightly lower values of the MIRO extrema near the cyclotron resonance. As shown in Fig. 4, both \(\kappa_{sh}\) (squares) and \(\kappa_{sm}\) (triangles) monotonically decrease with \(n_e\). The decrease in \(\kappa_{sm}\) with \(n_e\) occurs solely due to the weakening of the inelastic contribution, given by the second term in Eq. (5). This weakening, in turn, owes to a superlinear increase in the momentum relaxation time \(\tau\), which wins over the slightly sublinear increase in \(\tau_{sh}\) [see Eq. (4)]. In the smooth disorder limit, characterized by \(\kappa_{sm}\), the decrease becomes larger due to the growing ratio of \(\tau/\tau_{sh}\) which enters the denominator of the displacement contribution [first term in Eq. (5)]. We thus conclude that regardless of the exact disorder characteristics, theoretical predictions are in contrast with the experimentally obtained \(\kappa\) (circles) which shows significant increase over the density range studied [53]. Our findings were confirmed by measurements using \(f = 39.5\) GHz in another sample which is the case for all modern high-mobility GaAs quantum wells. This decrease should occur because \(\tau^{-1}\) [51] is less sensitive to small-angle scattering than \(\tau^{-1}\) and because the characteristic scattering angle decreases with density.

As shown in Fig. 4, both \(\kappa_{sh}\) (squares) and \(\kappa_{sm}\) (triangles) monotonically decrease with \(n_e\). The decrease in \(\kappa_{sm}\) with \(n_e\) occurs solely due to the weakening of the inelastic contribution, given by the second term in Eq. (5). This weakening, in turn, owes to a superlinear increase in the momentum relaxation time \(\tau\), which wins over the slightly sublinear increase in \(\tau_{sh}\) [see Eq. (4)]. In the smooth disorder limit, characterized by \(\kappa_{sm}\), the decrease becomes larger due to the growing ratio of \(\tau/\tau_{sh}\) which enters the denominator of the displacement contribution [first term in Eq. (5)]. We thus conclude that regardless of the exact disorder characteristics, theoretical predictions are in contrast with the experimentally obtained \(\kappa\) (circles) which shows significant increase over the density range studied [53]. Our findings were confirmed by measurements using \(f = 39.5\) GHz in another sample which are discussed in Supplemental Material.

We next examine the effect of density on the positions of the MIRO extrema near the cyclotron resonance. As noted from the data in Fig. 1(b), these extrema move closer towards \(\epsilon = 1\) with increasing density. To examine this behavior quantitatively, we introduce a parameter \(\varphi = (\epsilon^+ - \epsilon^-)/2\), where \(\epsilon^- (\epsilon^+)\) is the position of

FIG. 3. (Color online) \(\tau_{sh}\) (circles), obtained from the fits (cf. Fig. 2), and \(\tau_{q0}\) (squares), calculated using Eq. (3) and Eq. (4), vs \(n_e\). The dotted lines are guides to the eye.

FIG. 4. (Color online) \(\kappa\) (circles), obtained from the fits (cf. Fig. 2), \(\kappa_{sh}\) (squares), and \(\kappa_{sm}\) (triangles), calculated using Eq. (5) in the sharp and smooth disorder limit, respectively, vs \(n_e\). The dotted lines are guides to the eye.

\[
\frac{1}{\tau_q} = \frac{1}{\tau_{q0}} + \frac{1}{\tau_{ee}}. \tag{3}
\]

Under the conditions of our experiment the electron-electron scattering rate is given by [17, 42, 43]

\[
\frac{h}{\tau_{ee}} = \frac{\pi k_B^2 T^2}{4E_F} \ln \frac{2h v_F/a_B}{\pi k_B T^\star}, \tag{4}
\]

where \(E_F\) is the Fermi energy and \(a_B \approx 11\) nm is the Bohr radius in GaAs. Using Eqs. (3) and (4) we compute \(\tau_{q0}\) and present the results (squares) in Fig. 3. The results show that the impurity-limited quantum lifetime decreases slightly with \(n_e\), in general agreement with Ref. 31. We note, however, that in our experiment most of this decrease takes place at densities below \(\approx 2 \times 10^{11}\) cm\(^{-2}\).

As already mentioned, Fig. 2 also reveals a significant increase of the intercept of the Dingle fits, given by \(\mathcal{A}_0\), with increasing \(n_e\). Since \(\mathcal{A}_0 \propto \eta\) [47], this increase reflects the increase in \(\eta\), provided that that the density dependence of \(\mathcal{P}\) is accurately described by Eq. (2). To quantify this increase we introduce a parameter \(\kappa = \mathcal{A}_0(n_e)/\mathcal{A}_0(n_l)\), where \(n_l = 1.41 \times 10^{11}\) cm\(^{-2}\) is the lowest density studied. As shown in Fig. 4, \(\kappa\) (circles) increases by a factor of about three over the density range studied. This finding is unexpected since, as we show next, one should anticipate a decrease of \(\eta\) with increasing \(n_e\).

The dimensionless scattering rate \(\eta\) is given by [18]

\[
\eta = \frac{\tau}{2\tau_{\star}} + \frac{2\tau_{\star}}{\tau} \tau, \tag{5}
\]

where the first (second) term represents displacement [48–50] (inelastic [17, 18]) contribution. Here, \(\tau_{\star} \approx 8.2\tau_{ee}\) [17] and \(\tau/2\tau_{\star}\) [51] can vary between \(\tau/2\tau_{\star} = 6(\tau/\tau_{q0} + 3)^{-1}\) (smooth disorder limit) and \(\tau/2\tau_{\star} = 3/2\) (sharp disorder limit) according to the mixed disorder model [18, 52]. As a result, the relative change in \(\eta\) (or \(\mathcal{A}_0\)) with \(n_e\) is expected to fall between \(\kappa_{sm}\) and \(\kappa_{sh}\), given by \(\eta(n_e)/\eta(n_l)\) evaluated in the smooth and the sharp disorder limit, respectively. On a qualitative level, the decrease of \(\tau/2\tau_{\star}\) with \(n_e\) can be expected whenever \(\tau_{q}/\tau_{\star} \ll 1\), i.e., when small angle scattering dominates, which is the case for all modern high-mobility GaAs quantum wells. This decrease should occur because \(\tau_{\star}^{-1}\) [51] is less sensitive to small-angle scattering than \(\tau^{-1}\) and because the characteristic scattering angle decreases with density.

As shown in Fig. 4, both \(\kappa_{sh}\) (squares) and \(\kappa_{sm}\) (triangles) monotonically decrease with \(n_e\). The decrease in \(\kappa_{sm}\) with \(n_e\) occurs solely due to the weakening of the inelastic contribution, given by the second term in Eq. (5). This weakening, in turn, owes to a superlinear increase in the momentum relaxation time \(\tau\), which wins over the slightly sublinear increase in \(\tau_{sh}\) [see Eq. (4)]. In the smooth disorder limit, characterized by \(\kappa_{sm}\), the decrease becomes larger due to the growing ratio of \(\tau/\tau_{sh}\) which enters the denominator of the displacement contribution [first term in Eq. (5)]. We thus conclude that regardless of the exact disorder characteristics, theoretical predictions are in contrast with the experimentally obtained \(\kappa\) (circles) which shows significant increase over the density range studied [53]. Our findings were confirmed by measurements using \(f = 39.5\) GHz in another sample which are discussed in Supplemental Material.

We next examine the effect of density on the positions of the MIRO extrema near the cyclotron resonance. As noted from the data in Fig. 1(b), these extrema move closer towards \(\epsilon = 1\) with increasing density. To examine this behavior quantitatively, we introduce a parameter \(\varphi = (\epsilon^+ - \epsilon^-)/2\), where \(\epsilon^- (\epsilon^+)\) is the position of
the fundamental minimum (maximum). We then present obtained \( \varphi \) (circles) as a function of \( n_e \) and observe that it monotonically decreases with \( n_e \). Similar to \( \tau_{q0} \), the decrease is more pronounced at lower densities. Theory, however, predicts just the opposite behavior; as illustrated in Fig. 5(a), calculated values of \( \varphi_{sh} \) (squares) and \( \varphi_{sm} \) (triangles), representing sharp and smooth disorder limits, respectively, both increase with \( n_e \). The expected growth of \( \varphi_{sh} \approx \varphi_{sm} \) with \( n_e \) occurs, for the most part, due to the increase in \( \beta_{\omega} \), dominated by \( \tau_{em}^{-1} \propto n_e \), which controls the sharpness of \( \mathcal{P} \) near \( \epsilon = 1 \). Indeed, as shown in Fig. 5(b), \( \mathcal{P}(\epsilon) \) is considerably sharper at \( n_e = 1.41 \times 10^{11} \text{ cm}^{-2} \) (solid line) than at \( n_e = 2.87 \times 10^{11} \text{ cm}^{-2} \) (dotted line).

It is known that the phase reduction can occur with increasing \( \mathcal{P} \) due to contributions from multiphoton processes \([20, 21]\). This scenario, however, can be ruled out since \( \mathcal{P} \), in fact, decreases with \( n_e \) near \( \epsilon = 1 \) within investigated density range. As shown in Fig. 5, \( \mathcal{P} \) at the fundamental MIRO extrema (filled and open circles) exhibits a slight overall decrease within the studied density range, similar to \( \mathcal{P} \) at \( \epsilon = 1 \) (solid line). At higher MIRO orders, \( \mathcal{P} \) monotonically increases, as illustrated by dashed and dotted lines computed for \( \epsilon = 2 \) and \( \epsilon = 3 \), respectively. This increase in \( \mathcal{P} \) occurs because \( \beta_{\omega} \) remains relatively small within studied density range.

One somewhat uncertain parameter is \( \epsilon_{\text{eff}} \) which affects \( \beta_{\omega} \) entering \( \mathcal{P} \) given by Eq. (2). Indeed, the expression we have used is generally valid only when the overall sample thickness greatly exceeds the radiation wavelength, a condition which is not satisfied for the microwave frequency used in our experiment. According to Ref. 19, a better approximation would be using \( \epsilon_{\text{eff}} = 1 \) which would increase the value of \( \beta_{\omega} \) by approximately a factor of 2. However, any increase in \( \beta_{\omega} \) would only weaken (or even reverse) the density dependence of \( \mathcal{P} \) and further increase the disagreement between theory and experiment, both in \( \eta \) and \( \varphi \).

Finally, we note that by applying the gate voltage we are not only changing the carrier density but also modify the confinement potential. Numerical simulations show that the 2DEG is pulled away from the top interface towards the center of the quantum well and becomes wider with increasing density. Whether or not such a change of the confinement plays any significant role in the unexpected enhancement of the MIRO amplitude is unclear at this point and is left for future studies. To investigate this possibility, it would be interesting to perform measurements in different structures, such as heterojunction-insulated gate field-effect transistor, in which confinement becomes stronger with increasing carrier density.

In summary, we have investigated the effect of the carrier density \( n_e \) on the MIRO amplitude in a high-mobility modulation-doped GaAs/AlGaAs quantum well equipped with \textit{in situ} back gate. Our main finding is a significant \textit{growth} of the MIRO amplitude with increasing density. Dingle analysis shows that this increase originates primarily from \( \mathcal{P} \eta \) entering Eq. (1) and not from a slight increase of \( \tau_q \). This finding is in conflict with theoretical expectations which predict a modest increase of \( \mathcal{P} \) and a \textit{decrease} of \( \eta \) with increasing density. We further find that the MIRO extrema near the cyclotron resonance move toward each other with increasing \( n_e \) whereas the theory predicts just the opposite behavior. These findings indicate that our understanding of microwave photocurrent is still lacking and needs further examination.

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Mobility values were estimated using standard Van der Pauw method. Zero-field resistance anisotropy ratio in our samples was found to be less than two.

The obtained value of α is in agreement with recent theoretical calculations [36] considering GaAs quantum wells similar to the one employed in our study.


ψ was computed using m∗ = 0.0652, 0.0638, and 0.0631m0 for m∗ = 1.41, 2.15 and 2.87×1011 cm−2, respectively [28].


P was calculated using Eq. (2) with ϵac = 1 V/cm and m∗ = 0.067m0 entering ϵ and τme.


Microwave power is not exactly known and A0 differs from η by a numerical factor.


The rate of scattering on angle θ can be expressed in terms of angular harmonics, τθ = τn, as τnθ = i∞ τnθ eιθ. In this notation, τ0θ = τ0, τ1θ = τ1, τ2θ = τ2, ... In the mixed disorder model [52], τnθ = τnθ0δn0 + τnθsinθ for n = 1, 2, 3, ..., where τnθ0 and τnθsinθ represent, respectively, scattering rates for background impurities and remote ionized donors (located at a distance d from the interface, χ = (2kd)−2).
[53] An alternative theoretical proposal [54] considering MIRO in terms of classical memory effects for the case of small-angle scattering predicts $\eta \propto \tau n_c^{1/2}$. We note, however, that the inelastic mechanism was not considered in Ref. 54 and that their choice of the disorder potential predicts $\tau_q, \tau \propto n_c^{-1/2}$ leading to density-independent $\eta$.