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HEATING OF TWO-DIMENSIONAL ELECTRON GAS AND LO PHONONS IN δ-DOPED GaAs BY FAR-INFRARED RADIATION

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A change in the conductivity of the two-dimensional electron gas in δ-doped GaAs samples due to pulsed far-infrared laser irradiation has been detected. It is shown that the observed positive photoconductivity is caused by heating of 2D electrons. The energy loss of hot electrons is accomplished by scattering with nonequilibrium longitudinal optical (LO) phonons. An expression for the effective emission frequency of LO phonons coupled with electrons and the thermal bath of acoustic phonons is derived and the dependence of the emission frequency on the lattice temperature is experimentally obtained in the range 77–300 K.

Keywords: A. quantum wells, A. semiconductors, D. electron–phonon interaction, D. phonons, E. light absorption and reflection.

1. INTRODUCTION

Hot electron effects in the two-dimensional electron gas (2DEG) of III–V compounds have been extensively studied in the last years. One of the remarkable phenomena, characterizing 2DEG heating, is the formation of nonequilibrium longitudinal optical (LO) phonons, which gain their energy from hot electrons during the first stage of the cooling process. The lack of equilibration is caused by the strong polar interaction of electrons with LO phonons going along with an effective electron heating in two-dimensional systems at high electron densities including highly-degenerate 2DEG [1]. Previous investigations have been mainly carried out with 2DEG of quantum wells or of heterostructures at lattice temperatures near 4 K subjecting the electron gas to a substantial overheating (several 100 K) by powerful femto- and picosecond radiation pulses (see for instance [2] and references therein). In this case, the transfer of the energy from nonequilibrium LO phonons to the thermal bath of acoustic phonons is accomplished by the spontaneous decay of LO phonons into acoustic phonons of zone boundary states which are unoccupied at low lattice temperatures. The corresponding cooling time of the nonequilibrium system constituted by degenerate 2DEG and LO phonons is practically independent of the lattice temperature.

For structures with quantum wells or heterostructures, the formation of 2DEG is always accompanied by perturbations of the phonon spectrum like the occurrence of interface modes, phonon confinement etc. This makes the analysis of cooling phenomena more difficult and impedes a comparison with results of investigations of hot electrons and LO phonons in bulk semiconductors. From this point of view, the 2DEG of δ-doped structures is more suitable to study heating effects as in this case the lattice translation invariance is effectively undisturbed. The reason for the reduced cooling rate at high electron densities is not fully understood as yet [1, 2]. In addition to previous studies it is important to investigate the electron energy relaxation as a function of the lattice temperature. Here we report on measurements at comparatively small heating of 2DEG carried out by pulsed far-infrared laser in the range 77–300 K lattice temperature.

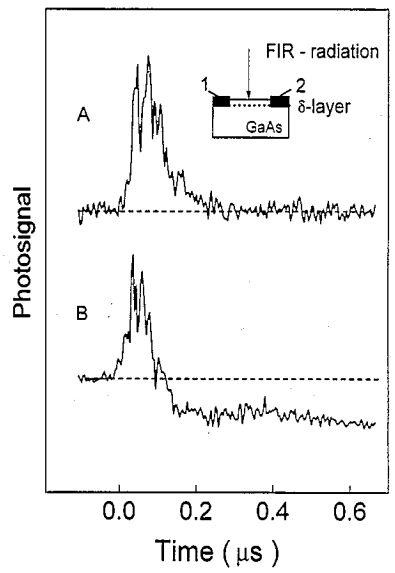

Fig. 1. The photosignal of the δ-doped structure at $T = 77$ K ($\lambda = 250$ μm). The positive sign of the signal corresponds to an increase in the 2DEG channel conductance under the action of the laser pulse. Pulse A corresponds to the "dark" condition when the sample was cooled in darkness, and pulse B was measured at steady-state illumination with visible light ("light-on" condition). In the last case, an additional long-time tail appears which occurs under the "light-on" condition only.

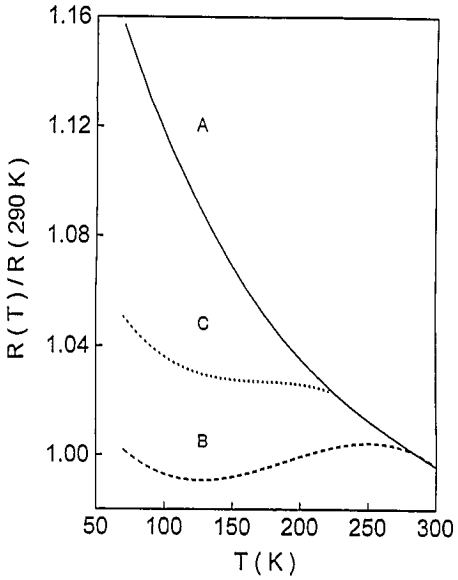

Fig. 2. The 2DEG channel resistance R as a function of the temperature T. Curve A was measured as the sample was cooled in darkness. Curve B corresponds to the "light-on" condition. Curve C was measured under the persistent photoconductivity condition, i.e. the sample was cooled in the dark from room temperature to $T = 77$ K and illuminated by visible light up to saturation of the change in the resistance. Then the $R(T)$ curve was measured in darkness.

2. SAMPLES AND EXPERIMENTAL TECHNIQUE

The investigated samples consist of MBE grown GaAs with one δ-layer, 5 mm long and 1 mm wide, at a distance of 20 nm from the semiconductor surface. The total wafer thickness is 300 μ m. The donor atoms are Si with a density of $6 \cdot 10^{12}$ cm^{-2}. The 2DEG density is $3 \cdot 10^{12}$ cm^{-2} due to the spatial redistribution of carriers between the surface states and the δ-layer. The investigations of the samples at temperature 4.2 K have shown, that tunneling-spectroscopy and magneto-transport data are in good agreement with the results of self-consistent calculations of the energy structure of two-dimensional subbands in the δ-doped layer of 5 nm width [3].

The used radiation source was a pulsed CH_3F molecular laser optically pumped by a TEA CO_2 laser providing $\leq$ 100 ns pulses at wavelength, λ, 250 μm. The maximum intensity of radiation in the sample was 1 MW/cm^2. For variation of the intensity, calibrated teflon attenuators were used. The intensity incident on the sample was simultaneously monitored by a fast photon drag detector. The measurements were carried out with the irradiation on the δ-layer side of the sample. Ohmic contacts to the δ-layer (1 and 2 in the insert of Fig. 1) were prepared on the surface of the sample and the sample was biased in series with a 50 Ohm load resistor. The sample resistance was $\geq$ 7 kOhm. The photoresponse, $\Delta\sigma/\sigma$, where σ is the conductivity of the sample, was determined from the voltage signal pulse.

3. EXPERIMENTAL RESULTS AND DISCUSSION

Figure 1 shows typical signal pulses. The sign of the fast signal corresponds to an increase in the 2DEG channel conductance due to far-infrared irradiation. This positive photoconductivity of the channel may be explained qualitatively by electron heating. To prove this, the temperature dependence of the channel resistance must be known. In Fig. 2, the measured resistance R as a function of the lattice temperature is shown for different conditions: (A) in the dark, (B) the sample permanently illuminated by visible light, (C) after illumination switched off (persistent photoconductivity). The origin of the persistent photoconductivity is likely to be the same as in Ref. [4]. In case of continuous illumination with visible light, a long-time

tail of opposite sign occurs in addition to the fast positive signal (Fig. 1, curve B). This will not be considered here because the fast channel signal related to 2DEG mobility change does not depend on the illumination conditions which affect mainly the 2DEG density.

Hall measurements have shown that the temperature dependence of resistance R in the dark is mainly caused by the change of 2DEG mobility . In order to analyze the heating effects in the 2DEG, the electron distribution in the two-dimensional subbands must be known. The shape of the δ-layer potential well and the spectrum of two-dimensional subbands were calculated in accordance with [3]. Then we obtained the value of the Fermi energy of 93 *meV* for the lowest and of 17 meV for the first excited subband. A rigorous treatment of hot 2D electrons is quite cumbersome taking into account the subband filling. As a first approach, we therefore have analyzed the measured photoresponse in a one-subband approximation. Effective parameters of the 2DEG were taken from Hall measurements.

As the mobility is determined by charged impurity scattering it depends on the electron temperature only. This allows us to use the dependence of R on T to determine the hot electron temperature T_e. The electron temperature approximation can be used in data handling. The presence of a Maxwell distribution of hot 2*D* electrons at a comparable electron density was proved experimentally by hot luminescence spectra in [5,6]. In addition, theoretical estimations of the electron-electron collision frequency give a time of 50 fs for the establishment of the Maxwell distribution of free electrons. This time constant is shorter than the electron - LO phonon relaxation time which is of the order of 150 fs [1].

Figure 3 shows the temperature dependence of the far-infrared response. The magnitude of the response drops by two orders of magnitude in a range where the temperature increases approximately by a factor of three. This behavior might be attributed to an exponential factor like $\exp(-\hbar\omega_0/k_BT_e)$ in the electron energy loss rate ($\hbar\omega_0 = 36.5$ meV is the energy of LO phonons in GaAs). As known, the electron energy loss rate due to scattering by LO phonons may be described by the term $A * \exp(-\hbar\omega_0/k_BT_e)$ irrespective of a possible heating of LO phonons (see for example [1]). The data of Fig. 3 show, however, that such an exponential term is not sufficient to describe the observed temperature dependence. In widely accepted theoretical expressions at the same time the pre-exponential factor A does not depend or shows only a weak dependence on the temperature (see for example Eqs. (3.10) and (3.13) in [1]).

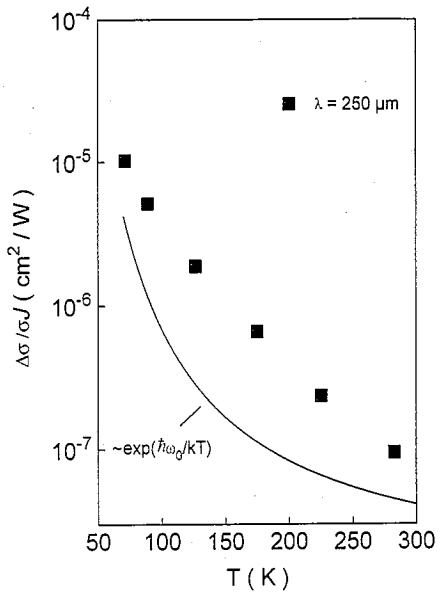

Fig. 3. Temperature dependence of the relative photoconductivity $\Delta\sigma/\sigma$ normalized to the intensity J. The solid line shows the $\exp(\hbar\omega_0/kT)$ behaviour.

As far as we know the explicit formulae for the electron energy loss by LO phonons which can be found in the literature were obtained for non-degenerate 2DEG only. Here we give a more general expression for this quantity which allows us to describe the experimental data given in Fig. 3 in the framework of a heating model.

4. THEORETICAL ANALYSIS OF ELECTRON-PHONON HEATING

In the electron temperature approximation, the set of equations for electron temperature T_e and for the nonequilibrium distribution function $N_\mathbf{f}$ of phonons is given by:

$$nc_e\frac{\partial T_e}{\partial t} = P_{\text{ext}} - P(T_e, N_\mathbf{f}), \tag{1}$$

$$\frac{dN_\mathbf{f}}{dt} = \frac{|c_\mathbf{f}|^2}{\hbar^2 V}\frac{\hbar\varkappa^2 f^2}{2\pi e^2}\left[-\text{Im}\varepsilon^{-1}(\mathbf{f}, \omega_\mathbf{f})\right] \times\left[N^0(T_e, \omega_\mathbf{f}) - N_\mathbf{f}\right] - \frac{N_\mathbf{f} - N_\mathbf{f}^0(T, \omega_\mathbf{f})}{\tau_\mathbf{f}(T)}. \tag{2}$$

Here n is the electron density, c_e is the heat capacity per electron, P_{ext} is the power per unit volume absorbed by electrons, $N_\mathbf{f}$ is the distribution function of phonons with wavevector $\mathbf{f}$, $N^0(T, \omega_\mathbf{f})$ is the Planck distribution function for temperature T, $|c_\mathbf{f}|^2$ is the square of the matrix element of electron-phonon interaction, V is the normalized volume, $\varkappa$ is the di-

electric constant of the lattice, $\varepsilon(\mathbf{f}, \omega_\mathbf{f})$ is the dielectric function of the electron gas, $\tau_\mathbf{f}(T)$ is the relaxation time of nonequilibrium phonons owing to the interaction with the thermal bath. It should be stressed that Eq. (2) is correct if only the phonon relaxation is due to creation-annihilation processes which are linear in the phonon occupation numbers $N_\mathbf{f}$ under the condition $N_\mathbf{f} << 1$. The system of Eqs. (1–2) is valid for phonons of any branch. In Ref. [7], it was applied to analyze the effect of the acoustic phonon heating on the dynamical energy relaxation time of the 3D electrons in *n*-InSb at quantizing magnetic fields. In that case, the relaxation time $\tau_\mathbf{f}(T)$ described the thermalization of acoustic phonons at the sample boundaries. Here we consider the decay of LO phonons into acoustic phonons.

The expression for the power $P(T_e, N_\mathbf{f})$ transferred per unit volume by electrons to phonons $N_\mathbf{f}$ can be obtained from Eq. (2) after multiplying the first term by $\hbar\omega_\mathbf{f}$ and taking the sum over all **f**. The result is:

$$P(T_e, N_\mathbf{f}) = \frac{1}{V}\sum_\mathbf{f} \hbar\omega_\mathbf{f}\nu_\mathbf{f}(T_e)\left[N^0(T_e, \omega_\mathbf{f}) - N_\mathbf{f}\right]. \tag{3}$$

Here we have introduced the emission frequency $\nu_\mathbf{f}(T_e)$ of phonons with wavevector **f** by the electron gas at temperature T_e:

$$\nu_\mathbf{f}(T_e) = \frac{|c_\mathbf{f}|^2}{\hbar}\frac{\varkappa^2 f^2}{2\pi e^2}\left[-\mathrm{Im}\varepsilon^{-1}(\mathbf{f}, \omega_\mathbf{f})\right]. \tag{4}$$

It should be noted that Eq. (3) is more general than the well- known expression of Kogan [8]. Here we did not make use of the assumption in [8] that the phonon subsystem is in thermal equilibrium. Thus the phonon distribution function $N_\mathbf{f}$ in Eq. (3) is arbitrary. The presence of $\mathrm{Im}\varepsilon^{-1}(\mathbf{f}, \omega_\mathbf{f})$ in Eqs. (1)-(4) allows for the screening of the electron–phonon interaction as well.

The duration of radiation pulses (~ 100 ns) used in this work are much longer than the characteristic cooling times of the electron–LO–phonon subsystem (a few ps [2]). Therefore, the radiation heating can be considered as stationary. Finding $N_\mathbf{f}$ from Eq. (2) for the steady state, substituting it in Eqs. (3) and (1) and disregarding the **f**-dependence of the LO phonon frequency ω_0, we obtain the equation for the electron temperature, which can be written as:

$$P_{\mathrm{ext}} = n\hbar\omega_0\nu(T_e, T)\left[N^0(T_e, \omega_0) - N^0(T, \omega_0)\right], \tag{5}$$

where we have introduced the effective emission frequency of LO phonons by electrons $\nu_e(T_e, T)$ which is given by:

$$\nu(T_e, T) = \frac{1}{nV}\sum_\mathbf{f}\frac{\nu_\mathbf{f}(T_e)}{1 + \tau_\mathbf{f}(T)\nu_\mathbf{f}(T_e)}. \tag{6}$$

All above expressions are valid for 3D electrons as well as for a 2DEG in single-band approximation if the proper matching of the dimension of the phonon wavevector **f** and of the expressions for $\varepsilon(\mathbf{f}, \omega_\mathbf{f})$ and $|c_\mathbf{f}|^2$ are made.

5. TEMPERATURE DEPENDENCE OF EFFECTIVE EMISSION FREQUENCY OF LO PHONONS

In the present experiments, the heating of the electrons is small ($\Delta T_e \approx 1 \div 20$ *K*). Therefore, in Eq. (5) $N^0(T, \omega_0)$ cannot be neglected in comparison to $N^0(T_e, \omega_0)$ as it is usually done. The values of $\nu_e(T_e, T)$ were obtained from the data by means of the following procedure. The electron temperature T_e was determined using the magnitude of the 2DEG photoresponse. The radiation-induced conductivity change $\Delta\sigma$ (Fig. 3) and $R(T)$ measured in the dark (Fig. 2) allow to calculate $N^0(T_e, \omega_0)$ in Eq. (5). After that the quantity

$$\nu_e^0(T) \equiv \nu_e(T_e, T)|_{T_e=T} \tag{7}$$

was calculated according to Eq. (5). The radiation power absorbed by the electrons, which is necessary for calculations, was evaluated according to the relation

$$P_{\mathrm{ext}} = \sigma(T)E^2,$$

where $\sigma(T)$ was taken from data in Fig. 2. The square of the effective electric field strength of the electromagnetic wave in the δ-layer was calculated from the measured intensity J of the incident radiation by the expression

$$E^2 = \frac{4R_{vac}}{(1 + \varkappa^{1/2})^2}J, \tag{8}$$

where $R_{vac} = 120\pi Ohm$ is the free space impedance. This expression does not take into account the possible interference of the radiation in the GaAs substrate.

The results of the calculations are given in Fig. 4. Note that the dependence of $\nu_e^0(T)$ on T in the investigated temperature range turned out to be strong enough that it must not be neglected. Since $\nu_e^0(T)$ at the low-temperature side was found to be of the order of the reciprocal cooling time, $1/\tau_a$, of nonequilibrium LO phonons in *GaAs* [1,2], [5,6], it is possible to assume that in this temperature range $\nu_e^0 \approx 1/\tau_a$ (see Eq. (6)). This can be seen from Eq. (6), assuming $\tau_\mathbf{f} \approx \tau_a = \mathrm{const}$ and $\nu_\mathbf{f}\tau_\mathbf{f} >> 1$ for all **f** in the range of essential contributions to the integral (the same analysis was done in [7]). Because $|\mathbf{f}|$ of LO phonons heated by electrons is small compared to wavevectors of acoustic phonons [9], the dependence of $\tau_\mathbf{f}$ on **f** can

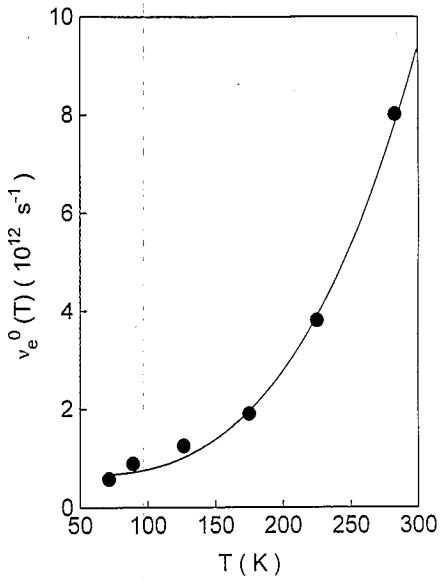

Fig. 4. The dependence of the effective emission frequency of LO phonons by electrons on the lattice temperature. The solid line is plotted in accordance to the Eqs. (10) and (11).

be neglected. Therefore it is possible to find a numerical approximation for the measured temperature dependence of $\nu_e^0(T)$ (Fig. 4). The temperature dependence of the relaxation time $\tau_a(T)$, which describes the contribution due to the temperature dependence of the occupation numbers of acoustic phonons, has been taken into account according to [9, 10]

$$\tau_a(T) = \tau_{ao} \tanh(\hbar\omega_0/4k_BT). \quad (9)$$

It turned out that the data given in Fig. 4 can be fitted well by

$$\nu_e^0(T) = \nu_{e0} \coth(\hbar\omega_0/4k_BT) P_2(k_BT/\hbar\omega_0), \quad (10)$$

where ν_{e0}=1.05·10^{12} s^{-1} and $P_2(x)$ is a second order polynomial of x, well described by

$$P_2(x) = 10x^2 - 4x + 1. \quad (11)$$

It is clear that Eq. (10) is asymptotically correct in the low-temperature range yielding $\nu_e^0(0) = \nu_{e0}$ in the limit $T \to 0$. At the same time, the dependence of $\nu_e^0(T)$ on T, following from Eqs. (10) and (11), does not agree with Eq. (9) in the high temperature limit. This disagreement may indicate that, as the lattice temperature grows, either the condition of strong LO phonon heating is violated or the theoretical derivation [10] of Eq. (9) must be revised. In the first case, the inequality $\nu_f\tau_f >> 1$ becomes invalid and $\nu_e(T_e, T)$ depends on ν_f which describes the electron energy relaxation to equilibrium LO phonons. The finite Fermi degeneration of the 2DEG and the dependence of the screening of electron–phonon interaction on the temperature can also be essential for the understanding of the observed dependence of $\nu_e^0(T)$ on T. With regard to the second case, we can add that the electron screening of phonon-phonon interaction should be taken into account at high electron densities in polar and piezoelectric semiconductors such as GaAs.

Note that the lattice heat capacity of GaAs as a function of temperature undergoes the crossover from the low-temperature quantum behavior to the high-temperature classical one [11] in the same temperature range of the order of 50–100 K where $\nu_e^0(T)$ deviates from Eq. (9). However, in this temperature range the Eq. (2) should be fully valid, as it follows from theoretical considerations in [12], for LO phonon - acoustic phonon coupling in GaAs up to the Debye temperature. This results from the fact that in polar semiconductors such as GaAs, which do not have inversion symmetry, three-phonon processes of decay (creation) of one LO phonon into (from) two acoustic phonons should be predominant (see Secstions 6 and 20 of [12]).

In conclusion, we have observed a far infrared photoconductive response of a 2DEG in δ-doped layers in GaAs. The change of conductivity is caused by heating of free electrons whose energy loss is due to emission of nonequilibrium LO phonons. The effective emission frequency $\nu_e(T_e, T)$ as a function of lattice temperature was determined from the photoconductive signal at low electron heating. The observed intensity dependence of the signal showed that $\nu_e(T_e, T)$ depends only weakly on T_e at $T < 100$ K. This result is in good agreement with previous investigations carried out at low lattice temperatures ($\sim$ 4 K) where in fact $\nu_e(T_e, 0)$ had been measured.

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