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2001 J. Phys.: Condens. Matter 13 3717

(http://iopscience.iop.org/0953-8984/13/16/303)

INSTITUTE OF PHYSICS PUBLISHING JOURNAL OF PHYSICS: CONDENSED MATTER

J. Phys.: Condens. Matter **13** (2001) 3717–3726 www.iop.org/Journals/cm PII: S0953-8984(01)19631-1

Nonlinear transport and optical properties of terahertz-driven two-dimensional electron gases

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Received 30 November 2000, in final form 5 February 2001

Abstract

A simple theory is developed for studying interactions between electrons and *intense* laser fields in semiconductor-based two-dimensional systems. With this approach, the steady-state transport and optical coefficients can be calculated easily for an electron–photon–phonon system. The results can be related to experiments where terahertz free-electron lasers are employed as intense radiation sources.

1. Introduction

In recent years, there has been a rapid expansion in developing high-power, long-wavelength, and frequency-tunable laser sources such as terahertz (THz) free-electron lasers (FELs). The current generation of the FELs can provide linearly polarized laser radiation which can be pulsed and/or in a continuum wave. Since 1995, THz FELs have been successfully applied in scientific research into optoelectronic properties in different semiconductor devices [1–4]. Motivated by the development and application of these state-of-the-art laser technologies, in this paper I study how electrons interact with linearly polarized intense THz laser fields in two-dimensional semiconductor systems (2DSSs).

When a semiconductor-based two-dimensional electron gas (2DEG) is subjected to intense THz laser fields, the electronic subband, Fermi, plasmon, and phonon energies of the system are of the order of the THz photon energy and, most importantly, the rate of electronic scattering with impurities and phonons is comparable to the frequency of THz photons. From a theoretical point of view, the conventional high-frequency dynamical theories and the normal linear-response approaches may not be applicable in dealing with the current situation. We have to develop a theory in which the radiation field is no longer treated as a perturbation. Recently, Lei proposed a balance equation approach [5] for the case where an electromagnetic (EM) field is present. This approach was based on modifications to the momentum- and energy-balance equations derived in the absence of the EM field [6]. In this paper, I develop a more tractable and more transparent theoretical approach in dealing with a 2DEG driven by an intense radiation field. In section 2, I study the electron–photon–phonon interactions in a 2DEG system. From these results, in section 3 I derive the momentum- and energy-balance equations based on the Boltzmann equation. These balance equations can be used to calculate transport and optical coefficients for a 2DEG driven by an intense THz laser field. The results obtained from this study are discussed in section 4 and are summarized in section 5.

2. Electron–photon–phonon interactions in a THz-driven 2DEG

In most experimental work published [1–3] on using THz FELs in the investigation of 2DSSs, the measurements are carried out for the situation where the laser field is polarized linearly along the 2D plane of a 2DEG (taken along the x-direction). Thus, the EM potential does not couple directly to the confining potential of the 2DEG (taken along the z-direction) and the radiation field plays a role as a driving field. That is why the set-up is called a 'THz-driven 2DEG'. It should be noted that the photon energy of THz radiation ($\hbar\omega \sim$ meV) is much less than the energy gaps ($E_g \sim$ eV) among different bands and valleys in semiconductor materials such as GaAs and Si. Therefore, the effects of interband and intervalley transitions through corresponding optical mechanisms can be neglected when the radiation intensity is not extremely high. As a consequence, the electron–photon–phonon interaction becomes the limiting factor in determining the transport and optical properties in 2DSSs driven by intense THz fields.

For a THz-driven 2DEG, the time-dependent electron Schrödinger equation can be solved analytically [7] and the effect of the radiation field can be included exactly within the time-dependent electron wavefunction $\psi_{n,\boldsymbol{k}}(\boldsymbol{R}, t)$. Applying $\psi_{n,\boldsymbol{k}}(\boldsymbol{R}, t)$ to time-dependent perturbation theory in which only the electron–phonon interaction is treated as a perturbation [8], the rate of the first-order steady-state electronic transition induced by electron–photon–phonon scattering is obtained as

$$W_{n'n}(\boldsymbol{k}', \boldsymbol{k}) = \frac{2\pi}{\hbar} \sum_{q_z} \left[N_Q + \frac{1}{2} \mp \frac{1}{2} \right] |V_Q|^2 G_{n'n}(q_z) \delta_{\boldsymbol{k}',\boldsymbol{k}+\boldsymbol{q}} \times \sum_{M=-\infty}^{\infty} J_M^2(r_0 q_x) \delta[E_{n'}(\boldsymbol{k}') - E_n(\boldsymbol{k}) - M\hbar\omega \mp \hbar\omega_Q]. \tag{1}$$

Here, n is the electronic subband index, $\boldsymbol{k} = (k_x, k_y)$ is the electron wavevector along the 2D plane, $\boldsymbol{Q} = (\boldsymbol{q}, q_z) = (q_x, q_y, q_z)$ is the phonon wavevector, $\hbar\omega_Q$ is the phonon energy, $N_Q = (\mathrm{e}^{\hbar\omega_Q/k_B T} - 1)^{-1}$ is the phonon occupation number, V_Q is the electron–phonon interaction coefficient, $G_{n'n} = |\langle n'|\mathrm{e}^{\mathrm{i}q_z z}|n\rangle|^2$ is the form factor induced by electron–phonon scattering in a 2DEG system, $|n\rangle$ is the electron wavefunction along the growth or z-direction, $E_n(\boldsymbol{k}) = \varepsilon_n + \hbar^2 k^2/2m^*$ is the energy spectrum of the 2DEG with m^* being the effective electron mass and ε_n the nth electron subband energy, and the upper (lower) case refers to absorption (emission) of a phonon. Furthermore, in equation (1), M is an index for M-photon absorption ($M > 0$) or emission ($M < 0$) and $M = 0$ is the index for elastic optic scattering, $r_0 = eF_0/(m^*\omega^2)$ with a dimension of length, ω and F_0 are respectively the frequency and the electric field strength of the radiation field, and $J_M(x)$ is a Bessel function.

Equation (1) exhibits features specific to electron–photon–phonon interactions in a 2DEG system when the radiation field is polarized linearly along the 2D plane. In the presence of intense EM radiation, conducting electrons in the system can interact with the radiation field via channels for optical emission and absorption. For a THz-driven 2DEG, the emission and absorption of photons by electrons can be achieved via indirect optical channels mediated by electron–phonon scattering. On the other hand, the emission and absorption of phonons via electronic transition events are accompanied by the emission and absorption of photons including multiphoton processes. For the case of high-frequency ($\omega \gg 1$) and/or low-intensity ($F_0 \ll 1$) radiation, entailing $r_0 \to 0$, equation (1) becomes that obtained by using Fermi's golden rule in the absence of the radiation field. From equation (1), it should be noted that the presence of the linearly polarized EM field can break the symmetry of the sample geometry, which results in an anisotropic electronic transition rate characterized by the dependence of $r_0 q_x$ via a term $J_M^2(r_0 q_x)$. This implies that the change of electron wavevector or momentum (via electron–phonon scattering) along the direction in which the radiation field is polarized, q_x, plays an essential role in switching different optical channels for electron–photon–phonon scattering, through the term $J_M^2(r_0 q_x)$ in equation (1).

For polar semiconductors such as GaAs-based 2DSSs, the frequency of the phonon oscillation associated with longitudinal optic (LO) modes is at the THz level ($\omega_{LO}/2\pi \simeq$ 8.85 THz). This suggests that the electron–LO-phonon coupling is the principal channel for relaxation of THz-excited electrons in a GaAs-based 2DEG, due to its large energy transfer during a scattering event. For electron–LO-phonon scattering: (i) $\omega_Q \simeq \omega_{LO}$, the LO-phonon frequency in the long-wavelength range; (ii) $N_Q \simeq N_0 = (\mathrm{e}^{\hbar\omega_{LO}/k_B T} - 1)^{-1}$; and (iii) the coupling coefficient is given by the Fröhlich Hamiltonian: $|V_Q|^2 = 4\pi\alpha L_0(\hbar\omega_{LO})^2/Q^2$, where α is the electron–LO-phonon coupling constant, and $L_0 = (\hbar/2m^*\omega_{LO})^{1/2}$ is the polar radius. Hence, the rate of the transition induced by the electron–photon–LO-phonon interaction reads

$$W_{n'n}(\boldsymbol{k}', \boldsymbol{k}) = 4\pi\alpha L_0 \hbar\omega_{LO}^2 \left[N_0 + \frac{1}{2} \mp \frac{1}{2}\right] X_{n'n}(q)\delta_{\boldsymbol{k}',\boldsymbol{k}+\boldsymbol{q}} \times \sum_{M=-\infty}^{\infty} J_M^2(r_0 q_x)\delta[E_{n'}(\boldsymbol{k}') - E_n(\boldsymbol{k}) - M\hbar\omega \mp \hbar\omega_{LO}] \tag{2}$$

where

$$X_{n'n}(q) = \int_{-\infty}^{\infty} \mathrm{d}q_z\, G_{n'n}(q_z)/(q^2 + q_z^2).$$

3. Momentum- and energy-balance equations

Generally, having obtained the electronic transition rate, the transport and optical coefficients can be calculated using, e.g., the Boltzmann equation. However, it is very difficult to solve this integro-differential equation, especially in the presence of the electron–phonon interaction, which is normally an inelastic scattering process. Thus, the balance equation approach has become a powerful tool in studying nonlinear properties of the electron gas systems. With this approach, one can circumvent the difficulties to solve the Boltzmann equation directly and keep the main merits of the Boltzmann equation as a governing transport equation. We have successfully applied this approach to the problem of hot-electron transport in 2DSSs in strong d.c. [9] and a.c. fields [10]. In this paper I generalize the balance equation approach for the case of a THz-driven 2DEG.

In the present study, I limit myself to the investigation of stationary electronic properties. The steady-state Boltzmann equation for a 2DEG in the case of nondegenerate statistics can be written as

$$\frac{\boldsymbol{F}}{\hbar} \cdot \nabla_{\boldsymbol{k}} f_n(\boldsymbol{k}) = g_s \sum_{n',\boldsymbol{k}'} [f_{n'}(\boldsymbol{k}')W_{nn'}(\boldsymbol{k}, \boldsymbol{k}') - f_n(\boldsymbol{k})W_{n'n}(\boldsymbol{k}', \boldsymbol{k})] \tag{3}$$

where $f_n(\boldsymbol{k})$ is the steady-state momentum distribution function for an electron in a state $|n, \boldsymbol{k}\rangle$, $g_s = 2$ accounts for spin degeneracy, $W_{n'n}(\boldsymbol{k}', \boldsymbol{k})$ is the steady-state electronic transition rate for an electron scattered from state $|n, \boldsymbol{k}\rangle$ to state $|n', \boldsymbol{k}'\rangle$, and $\boldsymbol{F}$ is the external force acting on the electron. For case where the transport and optical properties are investigated by means of conventional transport measurements, i.e., an external d.c. probing field E_x (or current I_x) is applied along the x-direction and the current (or voltage) is measured along this direction, we have $\boldsymbol{F} = -e(E_x, 0, 0)$. It should be noted that in the present problem, the effect of the radiation field has already been included within the electronic transition rate and, hence, the radiation field no longer appears in the force term of the Boltzmann equation.

For a THz-driven 2DEG, the momentum- and energy-balance equations can be derived using the approaches developed in reference [10]. The longitudinal resistivity of a 2DEG is defined by $\rho_{xx} = m^*/(n_e e^2 \tau)$, where n_e is the total electron density of the system and the momentum-relaxation time τ can be calculated from the momentum-balance equation through

$$\frac{1}{\tau} = \frac{4\hbar^2}{m^* n_e} \sum_{n',n} \sum_{k',k} k_x ({k_x}' - k_x) W_{n'n}(\boldsymbol{k}', \boldsymbol{k}) \frac{\partial f(\varepsilon_n + E)}{\partial E}. \quad (4)$$

Here $E = \hbar^2 k^2/2m^*$. Furthermore, the electron-energy-loss rate is defined as $P = \sigma_{xx} E_x^2/2$ with $\sigma_{xx} = 1/\rho_{xx}$, which can be determined from the energy-balance equation through

$$P = 2 \sum_{n',n} \sum_{k',k} [E_{n'}(\boldsymbol{k}') - E_n(\boldsymbol{k})] W_{n'n}(\boldsymbol{k}', \boldsymbol{k}) f(\varepsilon_n + E). \quad (5)$$

Introducing the rate of the transition induced by electron–photon–LO-phonon scattering, equation (2), into the momentum- and energy-balance equations, equations (4) and (5), we have

$$\frac{1}{\tau} = -\frac{\alpha L_0 \omega_{LO}^2 \sqrt{m^*}}{\sqrt{2}\pi^3 n_e} \left[N_0 + \frac{1}{2} \mp \frac{1}{2} \right] \sum_{M=-\infty}^{\infty} \sum_{n',n} \int \frac{\mathrm{d}^2 \boldsymbol{q}}{\varepsilon_q} \frac{q_x^2}{q} (\varepsilon_{n'nM}^{\mp} + \varepsilon_q) X_{n'n}(q) J_M^2(r_0 q_x)$$
$$\times \int_0^\infty \frac{\mathrm{d}E}{\sqrt{E}} \frac{\partial}{\partial E} f[E + \varepsilon_n + (\varepsilon_{n'nM}^{\mp} + \varepsilon_q)^2/(4\varepsilon_q)] \quad (6)$$

and

$$P = \sum_{M=-\infty}^{\infty} (M\hbar\omega \pm \hbar\omega_{LO}) \lambda_M \quad (7a)$$

with the inverse of the energy-relaxation time for the Mth optical process

$$\lambda_M = \frac{\alpha L_0 \omega_{LO}^2 m^*}{2\pi^3 \hbar n_e} \left[N_0 + \frac{1}{2} \mp \frac{1}{2} \right] \sum_{n',n} \int \frac{\mathrm{d}^2 \boldsymbol{q}}{\sqrt{\varepsilon_q}} X_{n'n}(q) J_M^2(r_0 q_x)$$
$$\times \int_0^\infty \frac{\mathrm{d}E}{\sqrt{E}} f[E + \varepsilon_n + (\varepsilon_{n'nM}^{\mp} + \varepsilon_q)^2/(4\varepsilon_q)]. \quad (7b)$$

Here, $\varepsilon_{n'nM}^{\mp} = \varepsilon_{n'} - \varepsilon_n - M\hbar\omega \mp \hbar\omega_{LO}$ and $\varepsilon_q = \hbar^2 q^2/2m^*$.

For case of a weak probing field E_x, the rate of electron energy loss induced by E_x is very small and $P \sim E_x^2 \ll 1$. Thus,

$$P_{op} = \sum_M M\hbar\omega\lambda_M \simeq P_{LO} = \sum_M (\mp)\hbar\omega_{LO}\lambda_M \quad (8)$$

where P_{op} and P_{LO} are the rates of the electron energy loss induced respectively by optical and phonon scattering processes. Equation (8) reflects the fact that the energy which an electron gains from the radiation field via optical absorption and emission is balanced by the energy which an electron loses via emission and absorption of phonons. Furthermore, the energy gain in an optical process is the summation of the processes caused by phonon emission and absorption, whereas the energy loss in a phonon process is the energy difference between the phonon emission and absorption.

For an electron gas system with relatively low electron density and at relatively high temperature or high excitation level, the Maxwellian with an electron temperature T_e, i.e., $f(x) = c\mathrm{e}^{-\beta x}$, is the most popularly used statistical electron energy distribution function, where $\beta = 1/k_B T_e$ and c is a normalization factor determined by the condition of electron number conservation via $c = \pi\hbar^2 n_e \beta/m^* A$ with $A = \sum_n \mathrm{e}^{-\beta\varepsilon_n}$. Using this distribution function, the momentum- and energy-relaxation times can be calculated through

$$\frac{1}{\tau} = \omega_{LO}\frac{4\alpha(\beta\hbar\omega_{LO})^{3/2}}{\pi^{3/2}A}\left[N_0 + \frac{1}{2} \mp \frac{1}{2}\right]\sum_{M=-\infty}^{\infty}\sum_{n',n}\mathrm{e}^{-\beta\varepsilon_n} \times \int_0^\infty \mathrm{d}q\, \frac{\varepsilon^{\mp}_{n'nM} + \varepsilon_q}{\hbar\omega_{LO}} X_{n'n}(q)\mathrm{e}^{-\beta(\varepsilon^{\mp}_{n'nM}+\varepsilon_q)^2/(4\varepsilon_q)} \int_0^1 \frac{\mathrm{d}x\, x^2}{\sqrt{1-x^2}}\, J_M^2(r_0 q x) \tag{9}$$

$$\lambda_M = \omega_{LO}\frac{2\alpha(\beta\hbar\omega_{LO})^{1/2}}{\pi^{3/2}A}\left[N_0 + \frac{1}{2} \mp \frac{1}{2}\right]\sum_{n',n}\mathrm{e}^{-\beta\varepsilon_n} \times \int_0^\infty \mathrm{d}q\, X_{n'n}(q)\, \mathrm{e}^{-\beta(\varepsilon^{\mp}_{n'nM}+\varepsilon_q)^2/(4\varepsilon_q)} \int_0^1 \frac{\mathrm{d}x}{\sqrt{1-x^2}}\, J_M^2(r_0 q x). \tag{10}$$

By solving the energy-balance equation, equation (8), we can obtain the electron temperature T_e and, then, calculate the power induced by optical P_{op} and phonon P_{LO} scattering processes. Introducing T_e into the momentum-balance equation, we can determine the momentum-relaxation time and resistivity. It is interesting to note that the inclusion of the laser field within the electronic transition rate, as proposed in this paper, can simplify the numerical calculations considerably. When taking the EM as an a.c. driving force in the drift term of the Boltzmann equation [10], one has to solve self-consistently the momentum- and energy-balance equations in order to obtain T_e. This is very CPU-time consuming. Employing the theoretical approaches developed in this paper, we can determine T_e by solving only the energy-balance equation, which is much easier to do numerically and can save a considerable amount of CPU time.

4. Numerical results and discussion

In this paper, the numerical calculations are carried out for GaAs-based single-quantum-well systems. The material parameters for GaAs are taken as: (i) the effective electron mass $m^* = 0.0665m_e$ with m_e being the rest electron mass; (ii) the electron–LO-phonon coupling constant $\alpha = 0.068$; and (iii) the LO-phonon energy $\hbar\omega_{LO} = 36.6$ meV. For the sake of simplicity, the confinement potential of the single quantum well along the growth direction is modelled using the infinite-height square-well approximation under which the eigenvalue $\varepsilon_n = \hbar^2 n^2\pi^2/2m^* L^2$ and eigenfunction $|n\rangle$ are known analytically, where L is the width of the square well. Thus, we have [11]

$$X_{n'n}(q) = \frac{\pi L}{4}\left[\frac{1}{a^2+c^2} + \frac{1+\delta_{n'n}}{b^2+c^2} - \frac{\pi^4 n'^2 n^2 c[1-(-1)^{n'+n}\mathrm{e}^{-2c}]}{(a^2+c^2)^2(b^2+c^2)^2}\right] \tag{11}$$

where $a = \pi(n'+n)/2$, $b = \pi(n'-n)/2$, and $c = qL/2$. In the calculations, two subbands, $n = 1$ and 2, are included. Moreover, the optical channels for $M = 0, \pm 1, \pm 2, \ldots, \pm 25$ are included within the numerical calculations. Inclusion of more optical processes affects only the results for low-frequency and very high-intensity radiation[1].

[1] In order to obtain the relatively good numerical results, one should take $|M| \gg \gamma = r_0/L_0$ in the calculations. It is found that when $f \sim 1$ THz and $F_0 \sim 10$ kV cm^{-1} (i.e., $\gamma \sim 15$), the contributions from $|M| > 25$ channels to τ and T_e are much smaller than those from $|M| < 15$ channels. It is less possible to achieve a multiphoton process with a larger M because it requires a larger energy transfer.

The dependence of the electron temperature on the THz radiation frequency is shown in figure 1 for different radiation intensities at a fixed lattice temperature and a fixed width of quantum well. At low radiation frequencies, the electrons in the system are cooled down, i.e., $T_e < T$. In this case, the process of an electron losing the energy via LO-phonon emission is much more efficient than that of gaining the energy via photon absorption, due to the large energy transfer caused by LO-phonon scattering. As a result, the requirement $P_{op} \simeq P_{LO}$ may lead to lowering of T_e with the result that the LO-phonon absorption becomes stronger. The cooling of electrons through a similar mechanism was observed in nonlinear electron transport in strong d.c. fields [9] and in the presence of a weak magnetic field [12]. With increasing radiation frequency $f = \omega/2\pi$, T_e first increases then decreases. In this case, an electron can be accelerated quickly to reach an energy level $E \geqslant \hbar\omega_{LO}$. When the time of this process required is shorter than the relaxation time required to emit a phonon, the electron will be heated, i.e., $T_e > T$. From figure 1, we see that for a GaAs-based 2DEG, the heating of electrons occurs at about $f \sim 1$ THz when $F_0 \sim 10$ kV cm^{-1}. In this radiation regime, T_e increases with F_0 and the peak of T_e shifts slightly to the higher-frequency regime with increasing F_0. At high radiation frequencies such that $\omega \sim \omega_{LO}$, an electron can easily gain energy, via photon absorption, to reach $E \sim \hbar\omega_{LO}$ and lose energy via phonon emission. This two-step process is analogous to an elastic scattering event and, hence, in this case the electrons will not be heated, i.e., $T_e \sim T$.

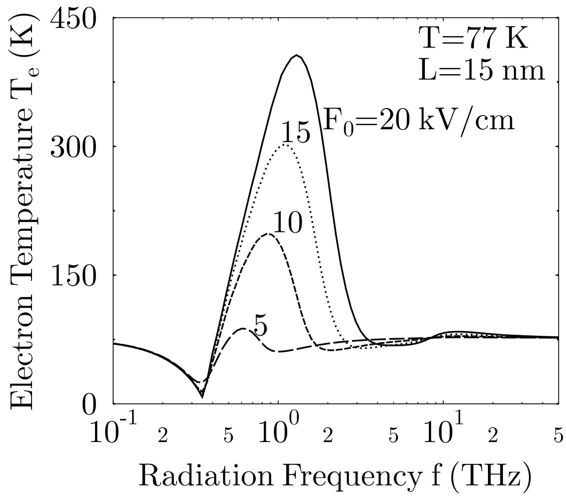

Figure 1. Electron temperature as a function of radiation frequency $f = \omega/2\pi$ for different radiation intensities F_0 at a fixed lattice temperature T and a fixed width of quantum well L.

The resistivity is shown in figure 2 as a function of radiation frequency for different radiation intensities at fixed lattice temperature, electron density, and width of quantum well. We see that an absorption peak can be observed at about $f \sim 1$ THz. With increasing radiation intensity, the absorption peak shifts to the higher-frequency regime. In 1995, Asmar *et al* [1] measured the frequency dependence of THz absorption for different 2DEG samples including GaAs-based heterojunctions and single quantum wells. They found an *unexpected* peak of resonant absorption at about $f \sim 1$ THz at moderate electron temperatures ($T_e \sim 80$ K) or radiation intensities. The results shown in figures 2 and 3 can be used to understand and interpret this important experimental finding. As can be seen from figure 3, at about $f \sim 1$ THz the strongest electron–photon–phonon scattering occurs via channels for optical

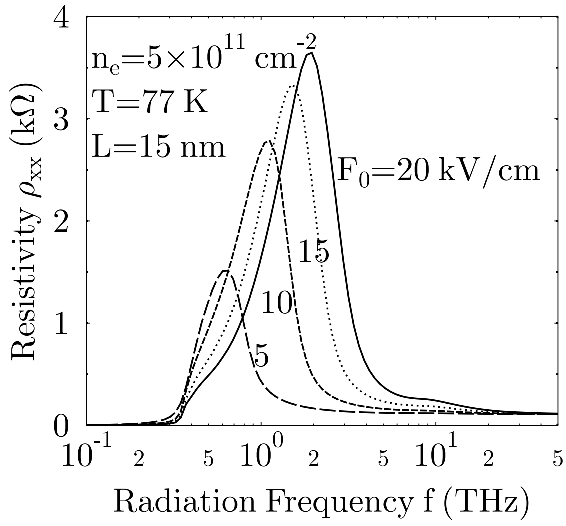

Figure 2. Resistivity as a function of radiation frequency for different radiation intensities at fixed lattice temperature, width of quantum well, and electron density n_e.

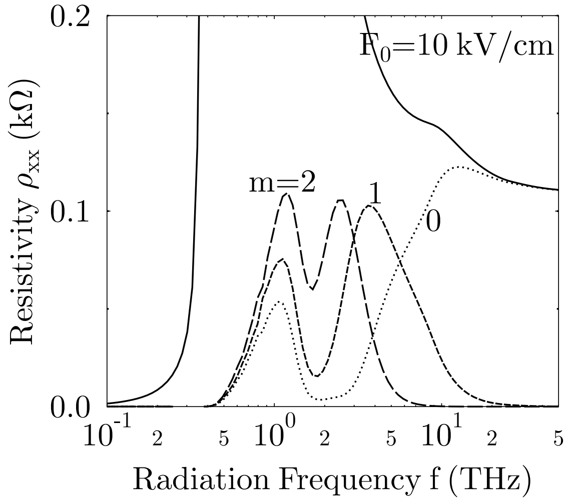

Figure 3. The contributions from different optical channels to the resistivity as functions of the radiation frequency at a fixed radiation intensity. $M > 0$ corresponds to a contribution from M-photon absorption and the solid curve shows the total resistivity. The other parameters are the same as for figure 2.

absorption including multiphoton absorption processes. When this happens, the electrons in the system are heated (see figure 1) and can gain energy from the radiation field and lose it by emission of LO phonons. This process opens up new channels for electronic transitions via resonant electron–photon–phonon scattering and, as a result, the resistivity increases and the optical absorption peak can be observed in the resistivity. From these theoretical results, we can conclude that the resonant electron–photon–LO-phonon coupling is the mechanism responsible for the resonant absorption of intense THz radiation observed experimentally.

In figure 4, the frequency dependence of the resistivity is shown at a fixed radiation intensity for different widths of the quantum well. We see that the absorption peak does not shift with varying quantum well width. This is in line with experimental findings [1] where the peaks of the resonant absorption were observed at about 1 THz for all samples such as heterojunctions and single quantum wells with different widths of the well layer. It should be noted in [1] that the experimental data were presented for a fixed electron temperature, because changing radiation frequency in FELs may result in a change of the radiation intensity. The results shown in figure 5 indicate that the resonant absorption of intense THz radiation can be measured not only at low lattice temperatures (as in reference [1]) but also at relatively high lattice temperatures through transport experiments.

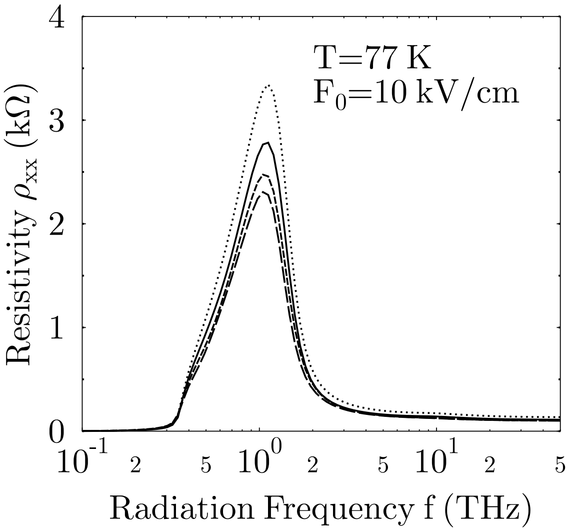

Figure 4. Resistivity as a function of radiation frequency at a fixed radiation intensity for different widths of the quantum well $L = 10$ nm (dotted curve), 15 nm (solid curve), 20 nm (dashed curve), and 25 nm (long-dashed curve).

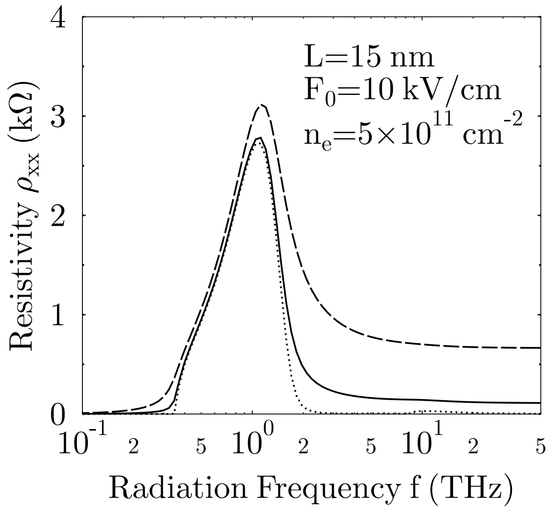

Figure 5. Resistivity as a function of radiation frequency at a fixed radiation intensity and a fixed quantum well width for different lattice temperatures $T = 10$ K (dashed curve), 77 K (solid curve), and 120 K (dotted curve).

The power of optical absorption per electron is plotted in figure 6 as a function of radiation frequency for different radiation intensities, and the contributions from different optical channels to the optical absorption at a fixed radiation intensity are shown in figure 7. From these results, we note that when the strongest electron–photon–phonon scattering occurs, the peak of the optical absorption can also be seen in the electron-energy-loss rate (EELR) at about $f \sim 1$ THz. Since the optical absorption coefficient is proportional to the EELR [5], we predict that the resonant absorption of the intense THz radiation by electrons in GaAs-based 2DEG systems, already seen in the transport experiments, will also be observable in optical measurements. I therefore hope that this can be verified experimentally.

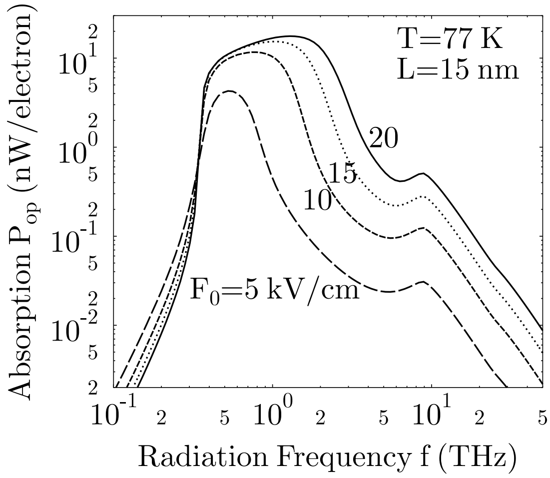

Figure 6. Power of optical absorption per electron as a function of radiation frequency for different radiation intensities at a fixed lattice temperature and a fixed width of quantum well.

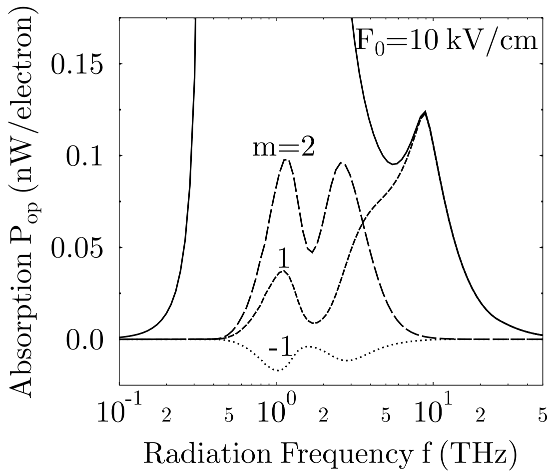

Figure 7. The contributions from different optical channels to the power of optical absorption per electron as functions of radiation frequency at a fixed radiation intensity. $M > 0$ ($M < 0$) corresponds to a contribution from M-photon absorption (emission) and the solid curve shows the total power of the optical absorption. The other parameters are the same as for figure 6.

5. Summary

In this paper, I have developed a simple theoretical approach for dealing with electron interactions with an intense laser field in an electron–photon–phonon system. Employing this approach, the electron–photon interaction can be included exactly and only the electron–phonon coupling is treated as a perturbation. The rate of the electronic transition induced by electron–photon–phonon interaction has been obtained for a THz-driven 2DEG. Introducing this electronic transition rate into the Boltzmann equation, I have derived the momentum- and energy-balance equations which can be used to calculate the transport and optical coefficients for the present system. Furthermore, I have carried out numerical calculations for electron temperature, resistivity, and power of optical absorption as functions of the THz radiation field, sample parameter, and experimental conditions. The results obtained from this study have been used to discuss and interpret those observed experimentally.

When a GaAs-based 2DEG is subjected to a laser field with a frequency $f \sim 1$ THz and an intensity $F_0 \sim 10$ kV cm^{-1} (or output power $P = 0.5\sqrt{\epsilon/\mu}|F_0|^2 \sim 100$ kW cm^{-2}), conditions such as $\omega \sim \omega_Q$ and $r_0 q \sim 1$ can be satisfied. As a consequence, the effects caused by electron–photon–phonon coupling can be observed. These radiation conditions have been realized by the current generation of THz FELs. Therefore, the simple theory developed in this paper can be applied to the investigation of THz-driven 2DEGs where THz FELs are used as intense THz radiation sources.

Acknowledgment

The author is a Research Fellow of the Australian Research Council.

References

[1] Asmar N G, Markelz A G, Gwinn E G, Černe J, Sherwin M S, Campman K L, Hopkins P F and Gossard A C 1995 *Phys. Rev.* B **51** 18 041
[2] Asmar N G, Černe J, Markelz A G, Gwinn E G, Sherwin M S, Campman K L and Gossard A C 1996 *Appl. Phys. Lett.* **68** 829
[3] Koenraad P M, Lewis R A, Waumans L R C, Langerak C J G M, Xu W and Wolter J H 1998 *Physica* B **256–258** 268
[4] Murdin B N, Heiss W, Langerak C J G M, Lee S-C, Galbraith I, Strasser G, Gornik E, Helm M and Pidgeon C R 1997 *Phys. Rev.* B **55** 5171
[5] Lei X L 1998 *J. Appl. Phys.* **84** 1396
[6] Lei X L and Ting C S 1985 *Phys. Rev.* B **32** 1112
[7] Xu W 1997 *Semicond. Sci. Technol.* **12** 1559
[8] Xu W 1998 *J. Phys.: Condens. Matter* **10** 6105
[9] Xu W, Peeters F M and Devreese J T 1991 *Phys. Rev.* B **43** 14 134
[10] Xu W and Zhang C 1997 *Phys. Rev.* B **55** 5259
[11] Xu W 1995 *Phys. Rev.* B **51** 13 294
[12] Xu W, Peeters F M and Devreese J T 1992 *Phys. Rev.* B **46** 7571