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Hot Electrons in Large-Bandgap Bulk Semiconductors and Quantum Wells

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Abstract

The analytical theory of hot electrons interacting with lattice vibrations, impurities, and interface roughness in quantum wells is developed. We have obtained new distribution functions which describe all the kinetic properties of the non-equilibrium electron gas. As a specific example we present the electric field dependence of the electron mobility in GaN-based quantum wells and in bulk GaN. The relative importance of the different scattering mechanisms is analysed in detail.

1. Introduction and formulation of the problem

The direct large-bandgap semiconductors like GaN are ideally suited for use in both electronic and optoelectronic devices as well as for combined applications. The desire to improve device performance has fuelled theoretical and experimental investigation of the optical and transport properties of these materials. Several GaN-based heterostructures have been grown recently and, with improving techniques, good quality GaN/GaAlN quantum wells (QW) will, no doubt, be produced in which the scattering of the electrons by the lattice vibrations is the dominant mechanism. A particularly important feature of these structures is doping which is a necessary stage in the sample preparation process in order to optimise carrier density. It is this feature that makes the study of the interactions of electrons with phonons and impurities in bulk GaN and GaN-based quantum wells of considerable importance.

Previous investigations of electron scattering phenomena in GaN were performed mostly for the equilibrium state of the electron system. Such a type of investigation gives a clear insight into the underlying physics and sheds light on the relative significance of different scattering mechanisms. It does lead to the prediction of the effect of lattice temperature, impurity doping, electron confinement, etc., on the linear electron transport in the systems in question. The linear response approximation with respect to

an external electric field greatly simplifies the problem because the appropriate electron energy distribution is then given by the equilibrium Maxwell-Boltzmann or Fermi-Dirac distribution function. The linear regime is thus applicable strictly only to the case of weak electric fields. On the other hand, the advent of high field electronic devices has necessitated the construction of an adequate theoretical framework for the non-equilibrium carrier dynamics of such devices. There are some purely numerical studies employing Monte Carlo techniques to model the non-equilibrium processes in the electron system of GaN. Such studies, by their nature, do not provide insight into the underlying basic physics. The main reason for resorting to numerical analysis stems from the difficulties of solving analytically the Boltzmann kinetic equation to obtain the electron distribution function appropriate at high electric fields. It should be noted that these difficulties occur for electron gases with high as well as low concentrations. At high electron concentrations one often uses the electron temperature approximation to solve the Boltzmann equation. But this approximation is not universally valid in the presence of polar optical scattering. The corresponding analysis for electrons in bulk materials [1] and in quantum wells [2] reveals very severe restrictions on the applicability of the electron temperature approximation. At low electron concentrations, on the other hand, the problem of finding the non-equilibrium distribution function can be tackled analytically in bulk materials for some practically important cases of electron transport [3]. This problem is much more complicated for electrons in low-dimensional systems with a quantized energy spectrum, especially if a few subbands are involved and it is necessary to solve a system of coupled Boltzmann equations [4].

In this paper we present the first theoretical analytical study of transport phenomena in bulk GaN as well as in GaN-based quantum wells under non-equilibrium conditions. The important feature of our theory is that it explicitly takes account of the many-subband electron population. We also assume that the electron average energy $\bar{\varepsilon}$ and the lattice temperature T satisfy the conditions, $(k_0 T, \bar{\varepsilon}) < \hbar\omega_0$, where $\hbar\omega_0$ is the optical phonon energy. These conditions impose an upper limit on both the external electric field E and the lattice temperature T. Since for GaN the optical phonon energy is large, $\hbar\omega_0 \approx 1100$ K, the condition can be satisfied within a wide range of actual E and T. Our recent investigation [5,6] of the electric field dependence of carrier mobility in low-dimensional structures has shown a non-monotonous behaviour and the dependence is characterised by a change in gradient when the scattering by the deformation acoustic (DA) phonons is more important than that by the piezoacoustic (PA) phonons. Here we take into account besides the DA and PA phonons also the interaction with polar optical (PO) phonons, background (BI) and remote impurities (RI), and with the QW interface roughness (IR). In our analysis the electron gas is assumed to be non-degenerate. This condition is not consistent with the electron temperature approximation. On going beyond the electron temperature approximation, we have obtained a set of new distribution functions for non-equilibrium electrons in GaN-based quantum wells with many electron subbands included. Our theory predicts novel non-linear regimes of transport phenomena in quantum wells and it provides a self-consistent transition from the two-dimensional (2D) regime to the three-dimensional (3D) regime of electron transport with increasing

electric field accompanied by the occupation of increasingly large number of subbands by the electrons.

2. The Boltzmann kinetic equation and the momentum relaxation times of 2D electrons

We consider here a 2D electron gas confined in a square infinitely deep AlN/GaN QW of thickness L subject to an external electric field $\vec{E}$ parallel to the QW layers. The energy spectrum of the electron with in-plane wavevector $\vec{k}_\parallel$ in subband n is $\varepsilon_n(\vec{k}_\parallel) = \hbar^2 k_\parallel^2/2m^* + W_0 n^2$ with $n = 1, 2, \ldots$ and $W_0 = \pi^2\hbar/2m^*L^2$, m^* is the electron effective mass. In order to calculate the kinetic coefficients of the non-equilibrium electron gas it is necessary to know the electron distribution function $F_n(\vec{k}_\parallel)$ which is governed by the Botzmann kinetic equation

$$-\frac{e\vec{E}}{\hbar}\frac{dF_n(\vec{k}_\parallel)}{d\vec{k}_\parallel} = \hat{I}F_n(\vec{k}_\parallel), \tag{1}$$

where $\hat{I}F_n(\vec{k}_\parallel)$ is a scattering operator which describes the interaction between electrons and the scatterers.

For interaction with the bulk-like phonons with wave vector $\vec{q} = (\vec{q}_\parallel, \vec{k}_\perp)$ and frequency $\omega_\lambda(\vec{q})$ we have

$$\begin{aligned}
&\hat{I}F_n(\vec{k}_\parallel) = \sum_{i=1}^{3}\hat{I}F_n(\vec{k}_\parallel) = \sum_{i=1}^{3}\sum_{n'}\sum_{\vec{q}} \mathrm{w_i}(\vec{q})G^2_{nn'}(q_\perp) \times \\
&\Big\{\Big[F_{n'}(\vec{k}_\parallel + \vec{q}_\parallel)(N_{\vec{q}}+1) - F_n(\vec{k}_\parallel)N_{\vec{q}}\Big]\,\delta\left(\varepsilon_{n'}(\vec{k}_\parallel + \vec{q}_\parallel) - \varepsilon_n(\vec{k}_\parallel) - \hbar\omega_\lambda(\vec{q})\right) + \\
&\Big[F_{n'}(\vec{k}_\parallel - \vec{q}_\parallel)N_{\vec{q}} - F_n(\vec{k}_\parallel)(N_{\vec{q}}-1)\Big]\,\delta\left(\varepsilon_{n'}(\vec{k}_\parallel - \vec{q}_\parallel) - \varepsilon_n(\vec{k}_\parallel) + \hbar\omega_\lambda(\vec{q})\right)\Big\}
\end{aligned} \tag{2}$$

Here $\mathrm{w_i}(\vec{q})$ is the scattering probability due to interaction with DA (i=1), PA (i=2), and PO (i=3) phonons, $N_{\vec{q}}$ is the phonon distribution function, the phonon frequency $\omega_\lambda(\vec{q}) = s_\lambda q$ for acoustic phonons (s_λ is the longitudinal ($\lambda = L$) or transverse ($\lambda = T$) acoustic velocity) and $\omega_\lambda(\vec{q}) = \omega_0$ for optical phonons, $G^2_{nn'}(q_\perp) = \left|(2/L)\int_0^L e^{-iq_\perp z}\sin(n\pi z/L)\right.$ $\left.\sin(n'\pi z/L)dz\right|^2$ is a form-factor. The scattering probabilities $\mathrm{w_i}(\vec{q})$ are equal, respectively: $\mathrm{w}_1(\vec{q}) = (\pi E_a^2/\rho V_0 s_L)q$, $\mathrm{w}_2(\vec{q}) = (\pi e^2 h_{14}^2/\rho V_0 s_T)C_\lambda(\vec{e}_q)q^{-1}$, $\mathrm{w}_3(\vec{q}) = (4\pi e^2\omega_0/V_0)$ $\left(\varepsilon_\infty^{-1} - \varepsilon_0^{-1}\right)$, where E_a is the deformation potential constant, h_{14} is the piezoelectric constant, $\varepsilon_0, \varepsilon_\infty$ are the low and high frequency dielectric constants, ρ is the material density, V_0 is the volume of the QW and $\vec{e}_q = \vec{q}/q$. The scattering probability $\mathrm{w}_2(\vec{q})$ by PA phonons depends in general on the orientation of $\vec{q}$ with respect to the crystal axes [7] through the factor $C_\lambda(\vec{e}_q)$ and on the crystal structure of GaN (zinc-blende or wurtzite). As a result $\mathrm{w}_2(\vec{q})$ for 2D electrons in a QW depends on the orientation of the interfaces of the quantum well with respect to the crystal axes. Here we use the PA-isotropy model in

which it is assumed that for a cubic crystal $C_L(\vec{q}) = 0$ and $C_T(\vec{q}) = 1$. Thus our results are more relevant to zinc-blende than the wurtzite structure of GaN.

For interaction with background impurities (BI) (i=4), remote impurities (RI) (i=5), and QW interface roughness (IR) (i=6), all of which are elastic, the scattering operator is given by

$$\hat{I}_i F_n^-(\vec{k}_\parallel) = \sum_{n'} \sum_{\vec{q}_\parallel} \left[\mathrm{w_i}(\vec{q}_\parallel) G_{nn'}^{(i)^2}(\vec{q}_\parallel) \left[F_{n'}^-(\vec{k}_\parallel + \vec{q}_\parallel) - F_n^-(\vec{k}_\parallel) \right] \right] \delta\left(\varepsilon_{n'}(\vec{k}_\parallel + \vec{q}_\parallel) - \varepsilon_n(\vec{k}_\parallel) \right), \tag{3}$$

where $\mathrm{w}_{4,5}(\vec{q}_\parallel) = \left(4\pi^2 n_s e^4 / A_0 \varepsilon_0^2\right) \vec{q}_\parallel^{-2}$, $\mathrm{w}_6(\vec{q}_\parallel) = \left(\pi^5 \hbar^4 \Delta^2 \Lambda^2 / A_0 m^* L^6\right) e^{-\frac{\vec{q}_\parallel^2 \Lambda^2}{4}}$, and the form-factors are $G_{nn'}^{(4,5)2}(\vec{q}_\parallel) = \left| \frac{2}{L} \int_0^L e^{-q_\parallel |z - z_d|} \sin(n\pi z/L) \sin(n'\pi z/L) dz \right|^2$, $G_{nn'}^{(6)^2} = 1$. Here n_s is a sheet impurity density, A_0 is the interface area; Δ and Λ are respectively the average height and correlation length of the interface roughness fluctuations and z_d is a position of the remote ($|z_d| > L/2$) impurities. For background impurities z_d is integrated over space.

In order to solve Eq.(1) we present $F_n(\vec{k}_\parallel)$ as the sum of a symmetric $F_n^+(\varepsilon_n)$ and an antisymmetric $F_n^-(\vec{k}_\parallel)$ function, $F_n(\vec{k}_\parallel) = F_n^+(\varepsilon_n) + F_n^-(\vec{k}_\parallel)$. The action of the scattering operator on the antisymmetric part of the distribution function is described by the electron momentum relaxation time, which is usually introduced by means of the equation

$$\hat{I}_i F_n^-(\vec{k}_\parallel) = -\frac{1}{\tau_n^{(i)}(\varepsilon_n)} F_n^-(\vec{k}_\parallel) \tag{4}$$

In the case of the 2D electron gas the problem is that in general the antisymmetric operator is described by a set of the relaxation times which correspond to intersubband and intrasubband scattering. This results in a coupling between different subbands in Eq.(1) which, in fact, becomes a system of the integro-differential equations. Our analysis of $\hat{I}_i F_n^-(\vec{k}_\parallel)$ in Eqs.(2), (3) reveals the following important features which drastically simplify following solution of Eq.(1):

(i). *Acoustic phonon scattering.* If the lattice temperature T is not very low, $k_0 T > \sqrt{8 m^* s_L^2 W_0}$, then one can treat the electron-phonon interaction as quasielastic and the antisymmetric operator can be written as

$$\hat{I}_i F_n^-(\vec{k}_\parallel) = -\sum_{i=1,2} \sum_{n'} \left(\Gamma_i^- / k_\parallel \right) \left[F_n^-(\vec{k}_\parallel) A_{nn'}^{(i)}(k_\parallel) - F_{n'}^-(\vec{k}_\parallel) B_{nn'}^{(i)}(k_\parallel) \right], \tag{5}$$

where

$$A_{nn'}^{(i)}(k_\parallel) = \int \frac{dq_\parallel}{\sqrt{1 - \varphi_{nn'}^2(q_\parallel)}} \int_{-\infty}^{+\infty} \frac{G_{nn'}^2(q_\perp) dq_\perp}{[q_\parallel^2 + q_\perp^2]^{(i-1)}}, \tag{6}$$

$$B^{(i)}_{nn'}(k_\parallel) = \int \frac{[1-(q_\parallel/k_\parallel)\varphi_{nn'}(q_\parallel)]dq_\parallel}{\sqrt{1-\varphi^2_{nn'}(q_\parallel)}} \int_{-\infty}^{+\infty} \frac{G^2_{nn'}(q_\perp)dq_\perp}{[q^2_\parallel + q^2_\perp]^{(i-1)}}, \tag{7}$$

and $\Gamma_1 = \hbar/2\pi m^* \lambda_a$, $\Gamma_2 = k_0 T/\pi\hbar s_T \tau_p$. Here $\lambda_a = \pi\hbar^4 \rho s^2_L/m^{*2} E^2_a k_0 T$ is a bulk electron mean free path due to DA phonon scattering and $\tau_p = 2\pi\rho\hbar^2 s_T/m^* e^2 h^2_{14}$ is a characteristic scattering time for the bulk material due to PA phonon scattering. The domain of integration for $q_\parallel$ in Eqs. (6) and (7) is defined by the condition $|\varphi_{nn'}(q_\parallel)| \le 1$, where $\varphi_{nn'} q_\parallel = q_\parallel/2k_\parallel - (2k_\parallel/q_\parallel)(n^2 - n'^2)W_0/4\varepsilon_\parallel$ and $\varepsilon_\parallel = \hbar^2 k^2_\parallel/2m^*$ is the electron kinetic energy in subband n. For DA phonons we obtain $A^{(i)}_{nn'}(k_\parallel) = \frac{2\pi^2}{L} k_\parallel \left(1 + \frac{1}{2}\delta_{nn'}\right)$ and $B^{(i)}_{nn'}(k_\parallel) = 0$. For PA phonons a numerical estimation of the integrals in Eqs. (6), (7) shows that only the intrasubband terms ($n' = n$) are important. This behaviour can be traced to the q-dependence of the integrals in Eqs.(6) and (7). It is a consequence of the $1/q^2$ factor in these equations since intersubband scattering requires a larger phonon wave vector in comparison with intrasubband scattering. Thus for PA phonon scattering we may neglect the intersubband terms ($n \neq n'$) in Eq.(5).The remaining intrasubband contribution in Eq.(5) has the form $\hat{I}_i F^-_n(\vec{k}_\parallel) = -\left(\Gamma^-_2/k_\parallel\right) F^-_n(\vec{k}_\parallel)\left[A^{(2)}_{nn}(\vec{k}_\parallel) - B^{(2)}_{nn}(\vec{k}_\parallel)\right]$. The coefficients $A^{(2)}_{nn}(\vec{k}_\parallel)$ and $B^{(2)}_{nn}(\vec{k}_\parallel)$ depend very weakly on the value of n since the form-factor $G^2_{nn'}(q_\perp)$ has the main maximum at $q_\perp = 0$ for all n. At $\varepsilon_\parallel \ll W_0$ we obtain $\left[A^{(2)}_{nn}(k_\parallel) - B^{(2)}_{nn}(k_\parallel)\right] \approx 2\pi$. As $\varepsilon_\parallel$ increases the difference $\left[A^{(2)}_{nn}(k_\parallel) - B^{(2)}_{nn}(k_\parallel)\right]$ decreases slowly, but one can ignore this change for larger $\varepsilon_\parallel$ due to the factor $\Gamma^-_2/k_\parallel \propto 1/\sqrt{\varepsilon_\parallel}$ the PA phonon scattering dominates at $\varepsilon_\parallel \ll W_0$. At higher energy the DA phonon contribution in Eq. (5) is more important than the PA phonon contribution. Finally, the antisymmetric operator due to acoustic phonon scattering has the form in Eq.(4) with the following momentum relaxation times: $\frac{1}{\tau^{(1)}_n(\varepsilon_n)} = \frac{\pi\hbar}{m^* L \lambda_a}\left[Z(\varepsilon_n) + \frac{1}{2}\right]$ for DA phonons, where $Z(\varepsilon_n) = int\left[(\varepsilon_n/W_0)^{1/2}\right]$, and $\frac{1}{\tau^{(2)}_n(\varepsilon_n)} = \frac{1}{\tau_p}\frac{2k_0 T_0}{\sqrt{2m^* s^2_T \varepsilon_\parallel}}$ for PA phonons.

(ii). *Polar optical phonon scattering.* If the condition $k_0 T \ll \hbar\omega_0$ is satisfied then only the absorption of the PO phonons should be taken into account. When an electron absorbs a phonon, it emits it almost instantaneously since the emission to the absorption time ratio is $\frac{\tau_{em}}{\tau_{ab}} \approx \frac{N_0}{N_0+1} \approx e^{-\frac{\hbar\omega_0}{k_0 T}} \ll 1$. In this process the electron energy remains almost unchanged, but the momentum changes drastically. Calculation in Eq. (2) shows that the intrasubband PO phonon scattering dominates over intersubband scattering and the antisymmetric operator has the form in Eq. (4) with a momentum relaxation time which is almost independent of the electron energy: $\frac{1}{\tau^{(3)}_n(\varepsilon_n)} = 3\sqrt{\frac{W_0}{\hbar\omega_0}} N_0 \frac{1}{\tau_0}$

(iii). *Background and remote impurity scattering.* Integration in Eq.(3) shows that here again only intrasubband terms are important because for intersubband terms $q_\parallel$ is large and the form-factor is exponentially small. The momentum relaxation times are similar to [7,8]: $\frac{1}{\tau^{(4)}_n(\varepsilon_n)} = \frac{\pi e^4 N_I L}{\varepsilon^2_0 \hbar\varepsilon_\parallel}\left[\pi - \frac{1}{2k_\parallel L}\ln\left(\sqrt{1+b_0} - b_0\right)\right]$ for background impurities, where $b_0 = \varepsilon_0\varepsilon_\parallel/e^2 N_I^{1/3}$ and N_I is the bulk impurity density, and $\frac{1}{\tau^{(5)}_n(\varepsilon_n)}\frac{\pi e^4 N_s}{\varepsilon^2_0\hbar\varepsilon_\parallel}[I_0(4k_\parallel z_d) -$

$L_0(4k_\parallel z_d)]$ for remote impurity scattering, where $I_v(x)$ and $L_v(x)$ are modified Bessel and Struve functions, respectively.
(iiii). *Interface roughness scattering.* For this scattering the intersubband terms are equal to zero and the momentum relaxation time due to intrasubband scattering is calculated similar to [9] as $\frac{1}{\tau_n^{(6)}(\varepsilon_n)} = \frac{\pi^5\Delta^2\Lambda^2\hbar}{m^*L^6}e^{-k_\parallel^2\Lambda^2/2}\left[I_0(k_\parallel^2\Lambda^2/2) - I_1(k_\parallel^2\Lambda^2/2)\right]$.

Substituting the momentum relaxation time approximation from Eq. (4) in Eq. (1) we obtain a solution of the Boltzmann equation for the antisymmetric distribution function in the form $F_n^-\vec{k}_\parallel) = (e\hbar/m^*)\tau_n(\varepsilon_n)(\vec{k}_\parallel\vec{E})dF_n^+(\varepsilon_n)/d\varepsilon_n$, where the total momentum relaxation time including all scattering mechanisms is defined as $\tau_n(\varepsilon_n) = \left(\lambda_a/\sqrt{2W_0/m^*}\right)\chi_n(\varepsilon_n)$ with $\chi_n(\varepsilon) = \left[(z(\varepsilon)+0.5) + \left(\lambda_a/\sqrt{2W_0/m^*}\right)\right.$ $\left.\sum_{i=2}^{6}\left(\tau_n^{(i)}(\varepsilon)\right)^{-1}\right]^{-1}$.

3. The energy distribution function of 2D electrons

The electric field dependence of the electron kinetic coefficients is described by the non-equilibrium symmetric distribution function. In order to find this function from the Boltzmann equation it is necessary to calculate the symmetric part of the scattering operator using Eq. (2). In this equation only interaction with DA and PA phonons should be taken into account. Using an approach similar to that developed in previous work [4] we obtain

$$\hat{I}F_n^+(\varepsilon_n) = \sum_{i=1,2}\sum_{n'}\Gamma_i^+ W_0\frac{d}{d\varepsilon_n}\left\{\left[\left(1+\frac{1}{2}\delta_{nn'}\right)\left(\frac{\varepsilon_n}{W_0}\right)^{2-i} - \left(1-\frac{i}{2}\right)n^2\delta_{nn'}\right]\times\right.$$
$$\left.\left[F_{n'}^+(\varepsilon_n) + k_0T\frac{dF_{n'}^+(\varepsilon_n)}{d\varepsilon_n}\right]\right\} + \sum_{i=1,2}\sum_{n'}\left(\Gamma_i^-/k_\parallel\right)A_{nn'}^{(i)}(k_\parallel)\left[F_{n'}^+(\varepsilon_n) - F_n^+(\varepsilon_n)\right], \tag{8}$$

where $\Gamma_1^+ = \delta_L(2\pi^2/L)\Gamma_1^-$, $\Gamma_2^+ = \delta_T L\Gamma_2^-$, and $\delta_\lambda = 2m^*s_\lambda^2/k_0T$.

The first term in Eq.(8) describes the rate of change of the energy distribution function $F_n^+(\varepsilon_n)$ due to *quasielastic* intra- and intersubband scattering. The second term describes the relaxation of $F_n^+(\varepsilon_n)$ due to *elastic* intersubband scattering. The total energy of the electron $\varepsilon_n = \varepsilon_\parallel + W_0$ does not change if the scattering is elastic, but the kinetic energy changes drastically if the electron undergoes intersubband scattering even if the total energy in the initial and final states is the same.

The equation for the symmetric distribution function is greatly simplified by the presence of the second term in Eq.(8), which is larger than the first term by the inverse quasielasticity factor $\delta_\lambda^{-1} \gg 1$. Retaining only this term in the kinetic equation one obtains, as a first approximation, the solution $F_{n'}^+(\varepsilon_n) = F_n^+(\varepsilon) \equiv F^+(\varepsilon)$ i.e. the symmetric distribution function does not depend on the subband index n [4]. The equation for

$F^+(\varepsilon)$ follows from Eq.(1) after the substitution from Eq.(8) and summation over n. The summation over n is necessary in order to eliminate the dependence of the coefficients in the equation on the subband index n since the distribution function does not depend on n. The equation obtained has the form of a continuity equation which, after integration, gives the solution

$$F^+(x) = A_0 \exp\left[-\frac{W_0}{k_0 T}\int_1^x \frac{dx'}{1+\varepsilon_E^2\alpha(x')}\right], \tag{9}$$

where $\alpha(x) = \sum_{n=1}^{Z(x)} \left(x-n^2\right)\chi_n(x)\left\{Z(x)\left(Z(x)+0.5\right)\left[x-\left(Z(x)+1\right)/6+\gamma_p s_T^2/2s_L^2\right]\right\}^{-1}$, A_0 is the normalization factor, $x = \varepsilon/W_0$ and $\varepsilon_E = eE\lambda_a/2\sqrt{m^* s_L^2 W_0}$.

4. Mobility of 2D and 3D electrons in GaN

The distribution functions obtained describe all the kinetic properties of 2D non-equilibrium electrons in the QW and this is the central result of the paper. As a specific example, we have performed calculations of the electron mobility μ as a function of the electric field E for a GaN-based QW. For comparison we have also plotted the corresponding dependence for bulk GaN. We have obtained the following expression for the electron mobility of 2D electrons:

$$\mu(E) = -\frac{e\lambda_a L}{\pi\hbar}\sum_{n=1}^{\infty}\int_{n^2}^{+\infty}\left(x-n^2\right)\chi_n(x)\frac{dF^+(x)}{dx}\Bigg/\sum_{n=1}^{\infty}\int_{n^2}^{+\infty}F^+(x)dx, \tag{10}$$

We have used the following values of GaN parameters: $m^* = 0.15m_0$, $E_a = 10.1eV$, $h_{14} = 4.24\times10^7$ V/cm, $\varepsilon_0 = 9.5$, $\varepsilon_\infty = 5.4$, $\rho = 6.1$ g/cm^3, $S_L = 4.57\times10^5$ cm/s, $s_T \approx 0.5 s_L$, $\hbar\omega_0 = 92.8$ meV. To add some imperfections we have also used the parameters $N_I = 10^{16}$ cm^{-3} for BI scattering $N_s = 10^{11}$ cm^{-3} and $z_d = L + 500\mathring{A}$ for RI scattering and $\Lambda = 20\mathring{A}$ and $\Delta = 2\mathring{A}$ for IR scattering.

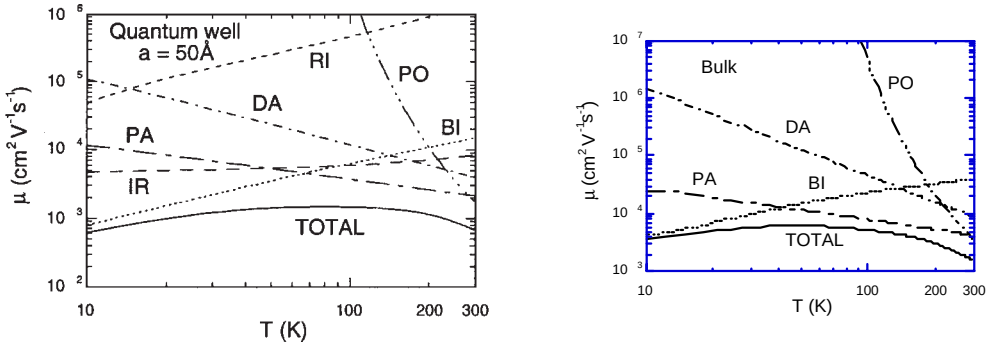

Figure 1. Low field mobility of electrons versus temperature for a GaN quantum well and bulk GaN.

First we show in Figure 1 the low-field mobility, $E \to 0$ as a function of temperature. This

figure clearly demonstrates the relative importance of different scattering mechanisms. Figure 2 shows the electric field dependence of the mobility $\mu(E)$ at different temperatures and well widths. It can be seen that temperature has a significant effect, not only on the magnitude of the mobility, but it changes also the style of the dependence. At low temperatures $\mu(E)$ is a non-monotonous function of E with a maximum which is caused by the competition between impurity, IR and PA scattering and DA phonon scattering. As T increases imperfections become less important and the PO phonon contribution to μ starts to dominate. Because the momentum relaxation time due to PO phonon scattering is almost independent on the electron energy, the mobility depends only weakly on electric field.

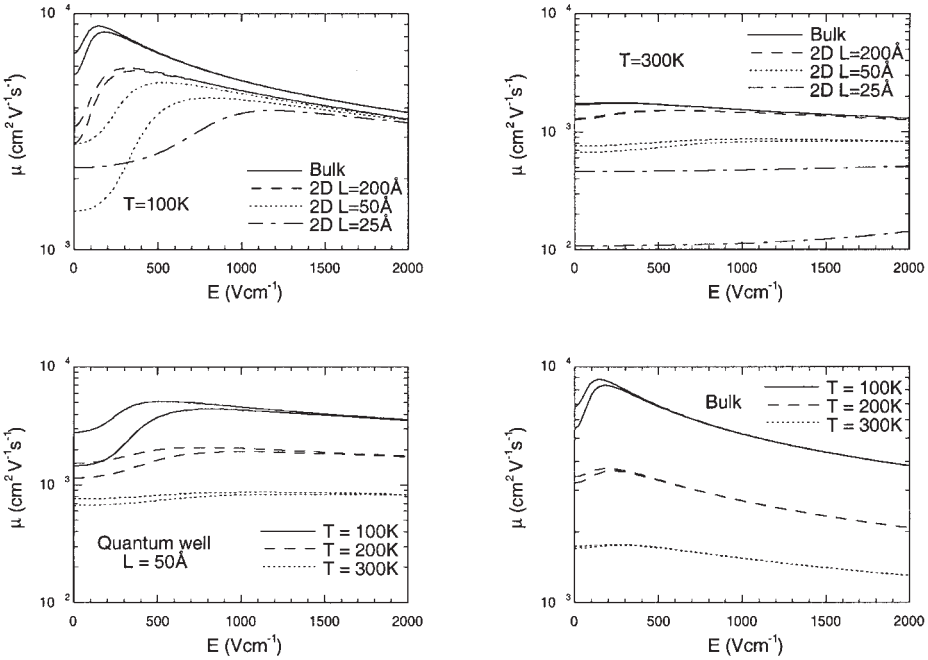

Figure 2. Field dependence of the mobility of electrons for different lattice temperatures and well widths. The lower curves for each case are for an imperfect sample including IR, RI and BI scattering as well as phonon scattering while the higher curves include only phonon scattering. The mobility for the imperfect 25$\mathring{A}$ well at T=100K is too small to appear.

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