Electron-energy-loss rates in $Al_xGa_{1-x}As/GaAs$ heterostructures at low temperatures

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We have measured the energy-loss rates of electrons in modulation-doped $Al_xGa_{1-x}As/GaAs$ heterojunctions at low temperatures. At the temperatures of this experiment, the energy loss is due to acoustic-phonon scattering, mostly through the deformation-potential interaction. We describe a theory of the energy-loss rate to acoustic phonons, discuss the effects of static screening, and determine a value of the deformation-potential constant by matching the theory to our experimental data. We show how the temperature dependence of the energy-loss rate can be used to determine the effectiveness of screening. In agreement with previous studies, we find that an anomalously large deformation-potential constant of 16.0 or 11.5 eV, for the screened or unscreened theory, respectively, is necessary to explain the data. However, in contrast with previous studies, we do not interpret this result as indicating some error in previous measurements of the deformation potential in bulk GaAs. Rather, we suggest that it results from some inadequacy in the model used to represent acoustic-phonon scattering in heterojunctions, or from some additional scattering mechanism not present in bulk material.

I. INTRODUCTION

Recent studies of acoustic-phonon scattering in $Al_xGa_{1-x}As/GaAs$ heterojunctions have led to controversial conclusions regarding acoustic deformation-potential scattering.[1-4] In a modulation-doped heterostructure the electronic wave function is largely confined to the GaAs, and we therefore expect the deformation-potential constant Z to be that of bulk GaAs. However, a number of studies of transport properties in $Al_xGa_{1-x}As/GaAs$ heterostructures have led to the suggestion that the deformation-potential constant is 11—16 eV,[1,5-7] considerably larger than the commonly accepted bulk value of 7—8 eV.[8,9] Because transition rates calculated in the deformation-potential formalism are proportional to Z^2, this is a significant discrepancy.

Most previous studies of the deformation-potential interaction in heterojunctions have inferred a value for Z from the temperature dependence of the mobility. It is known that the absolute mobility is limited mainly by ionized impurity scattering, even in samples with mobilities of the order of 10^6 $cm^2/V\,sec$, but it is not clear which scattering mechanism dominates the temperature dependence of the mobility. These studies assumed that, at temperatures below ~ 40 K, the temperature dependence of the mobility is due solely to acoustic-phonon scattering. However, it has been shown that the temperature dependence of ionized impurity scattering is significant in samples which are not highly degenerate.[4] Lin, Tsui, and Weimann have found that ionized impurity scattering dominates the temperature dependence of the mobility in samples with mobilities less than approximately 10^5 $cm^2/V\,sec$ at 4.2 K.[5] The effect of ionized impurity scattering is reduced in samples with high electron concentration and high mobility, but may still be important. Other scattering mechanisms, such as surface roughness or surface charge scattering, may also affect the mobility. These contributions to the mobility can mask the effect of phonons and complicate the determination of Z. This possibility is of special concern because of the discrepancy between the value of Z necessary to fit data from bulk material and the value necessary to fit data from heterojunctions.

A recent study by Hirakawa and Sakaki[7] has used the electronic-energy-loss rate at a lattice temperature of 4.2 K to determine Z. The energy-loss rate is unaffected by elastic scattering mechanisms, such as ionized impurity scattering, and the results are therefore more reliably interpreted. At temperatures below ~ 40 K, where polar optical-phonon scattering is negligible, acoustic phonons provide the only known inelastic scattering mechanism. They determined that a deformation potential constant of 11 eV was necessary to fit their energy loss data to a theory which included screening of the electron-phonon interaction.

In this study we measure the energy-loss rate at various lattice temperatures, present a theory of the energy-loss rate to acoustic phonons, including a discussion of screening effects, and suggest possible reasons for the discrepancy between the value of the deformation potential necessary to fit data from bulk and from heterostructures.

II. THEORY

The theory of scattering in quasi-two-dimensional systems was first developed to explain mobility and energy loss in the inversion layer formed in Si metal-oxide-semiconductor field-effect transistors (MOSFET's).[10-13] It has been shown that, at low temperatures, deviations from equipartition result in an energy-loss rate with a T_l dependence determined by the Q dependence of $|M(\mathbf{Q})|$, where T_l is the lattice temperature, Q is the phonon wave

vector, and $M(\mathbf{Q})$ is the scattering matrix element (including the screening factor, but not the phonon occupation number). With the simplifying assumption of strictly two-dimensional (2D) phonons, $|M(\mathbf{Q})|^2 \propto Q^p$ leads to $\langle \partial\varepsilon/\partial t \rangle \propto (T_e - T_l)T_l^{2+p}$, where $\langle \partial\varepsilon/\partial t \rangle$ is the average energy-loss rate per electron, and T_e is the electron temperature.[10] If we assume 3D phonons, we find $\langle \partial\varepsilon/\partial t \rangle \propto (T_e - T_l)T_l^{3+p}$. The dependence of the energy-loss rate on the Q dependence of the scattering matrix element provides an important method of distinguishing experimentally between different scattering mechanisms. In particular, as we will discuss later, the effect of screening can be deduced from the T_l dependence of the energy loss because the screening introduces an extra factor of Q^2 into $|M(\mathbf{Q})|^2$. These early studies of scattering assumed that the electrons were purely 2D. This leads to results which are quantitatively incorrect.

More recent work has assumed that the electrons are quasi-2D. That is, the electronic envelope wave function is given by

$$\Psi \propto \exp(i\mathbf{k}_{||}\cdot\mathbf{r}_{||})\phi(z) , \tag{1}$$

where the wave function is separable into plane waves parallel to the confining interface, and a narrowly localized function $\phi(z)$ perpendicular to the interface. At the temperatures and electron densities of our experiment, only the first subband is occupied, and all electrons have the same perpendicular wave function.

Price has shown that the use of quasi-2D wave functions leads to matrix elements which are proportional to the overlap integral[14]

$$I(Q_\perp) = \int \phi^2(z)\exp(iQ_\perp z)dz , \tag{2}$$

where $Q_\perp$ is the perpendicular component of the wave vector of the phonon with which the electron interacts. The overlap integral is 1 for small $Q_\perp$, and falls off to 0 as $Q_\perp$ becomes large. How fast it falls off depends on the electron wave function. For the self-consistent wave function used in this study, $I^2(Q_\perp)=0.9$ at $Q_\perp \simeq 6\times 10^5$ cm^{-1}, and decreases to 0.1 at $Q_\perp \simeq 4\times 10^6$ cm^{-1}. At low temperatures the results are less sensitive to the exact form of the wave function because the overlap integral is 1 for all phonon modes with significant occupation numbers. For a typical wave function in an $Al_xGa_{1-x}As/GaAs$ heterojunction, the lattice temperature must be below 1 K for the approximation $I^2(Q_\perp)=1$ to be accurate. At the temperatures of our experiments (1.4–9 K), the results are still sensitive to the specific wave function used.

The exact form of the wave function can have a strong effect on $I(Q_\perp)$ and therefore on the average scattering rates. In our study, the wave function perpendicular to the interface is calculated using a self-consistent interface potential, including exchange-correlation effects. We assume a conduction-band offset of 300 meV. For comparison, we have also used the approximate Fang-Howard variational wave function, with somewhat different results. We will discuss this later.

Price has suggested that screening of the deformation potential must be included.[15] Static screening in the random-phase approximation is accounted for by multiplying the unscreened matrix element by a screening factor, S, given by

$$S = \frac{Q_{||}}{Q_{||} + PH(Q_{||})} , \tag{3}$$

where

$$H(Q_{||}) = \int\int dz_1 dz_2 \phi^2(z_1)\phi^2(z_2)\exp(-Q_{||}|z_1 - z_2|) , \tag{4}$$

and

$$P = (e^2/\varepsilon_r\varepsilon_0)(2m^*/\hbar^2) . \tag{5}$$

For GaAs, $P \simeq 2\times 10^6$ cm^{-1}.

It is desirable to verify the accuracy of this screening theory experimentally. The theory of screening in an anisotropic system is complicated and requires the use of untested approximations. It is also unclear how accurately the static (instead of dynamic) screening theory describes the electron-phonon interaction.

In the limit of strong screening, the static screening factor reduces to

$$S = \frac{Q_{||}}{P} . \tag{6}$$

As noted earlier, the extra factor of Q changes the T_l dependence of the energy-loss rate at low lattice temperatures. (In the approximations that allow for an analytic solution, the parallel and perpendicular components of Q contribute to the power of the T dependence in the same way as the full 3D component.) We find that for $(T_e - T_l) < T_l$ the deformation-potential interaction leads to $\langle \partial\varepsilon/\partial t \rangle \propto (T_e - T_l)T_l^4$ if screening is not included, and to $\langle \partial\varepsilon/\partial t \rangle \propto (T_e - T_l)T_l^6$ if screening is included.[10,16] At high lattice temperatures, the energy loss is proportional to $(T_e - T_l)$ for both the screened and the unscreened interactions. Although these temperature dependences give a qualitative understanding of the energy-loss rate, they depend on approximations, and only apply over limited temperature ranges. The experiments reported in this study have been done in an intermediate-temperature range where there is no simple relationship between the energy-loss rate and the lattice temperature.

Current theories of scattering in $Al_xGa_{1-x}As/GaAs$ heterojunctions assume that the phonon modes are the same as those of bulk GaAs. This may be an oversimplification. It seems probable that long-wavelength phonons will be affected by the proximity of the surface of the crystal. Mendez, Price, and Heiblum justify the assumption of bulk 3D phonon modes by citing magnetophonon experiments which indicate that the energy of the polar optical-phonon modes is relatively unchanged,[1] but this energy is probably insensitive to the mode of the phonon, and therefore a poor indicator of any disturbance in the modes. Studies of the mobility in Si MOSFET inversion layers have shown that interface modes can be important in that system.[13] Interface phonons have been clearly observed in $Al_xGa_{1-x}As/GaAs$ superlattices,[17] but we know of no conclusive work on single heterointerfaces. We continue with the assumption of bulk 3D phonon

modes because there exists no definitive theory of interface modes, but we caution that the results should not be compared directly with those from bulk material.

In order to calculate the energy-loss rate it is necessary to know how the applied electric field affects the form of the electron distribution function. We assume that the electrons have a Fermi-Dirac distribution $f(\varepsilon)$ with an electron temperature T_e which is greater than the lattice temperature T_l.

It is convenient to calculate the average energy loss per electron by calculating the energy gained by the phonons from the electrons and dividing by the number of electrons which participate.[18] Using this approach, we write

$$\left\langle \frac{\partial \varepsilon}{\partial t} \right\rangle = \frac{-1}{N_e} \sum_{\mathbf{Q}} \hbar\omega_Q \frac{\partial N_Q}{\partial t} , \tag{7}$$

where N_e is the number of electrons participating, and $\hbar\omega_Q$ is the energy of a phonon mode. $\partial N_Q/\partial t$ may be written as

$$\begin{aligned} \frac{\partial N_Q}{\partial t} = 2\frac{2\pi}{\hbar} \sum_{\mathbf{k}} & I^2(Q_\perp) S^2(Q_{||}) \, | M(\mathbf{Q}) |^2 \\ & \times \delta(\varepsilon_{\mathbf{k}} + \hbar\omega_Q - \varepsilon_{\mathbf{k}+\mathbf{Q}_{||}}) \\ & \times \{ (N_Q + 1) f(\varepsilon_{\mathbf{k}+\mathbf{Q}_{||}}) [1 - f(\varepsilon_{\mathbf{k}})] \\ & - N_Q f(\varepsilon_{\mathbf{k}}) [1 - f(\varepsilon_{\mathbf{k}+\mathbf{Q}_{||}})] \} , \end{aligned} \tag{8}$$

where N_Q is the Planck distribution function, ω_Q is related to Q by the dispersion relation $\omega_Q = uQ = u(Q_\perp^2 + Q_{||}^2)^{1/2}$, and $\varepsilon_{\mathbf{x}} = (\hbar^2/2m^*) | \mathbf{x} |^2$. Here u is the sound velocity, either longitudinal or transverse, ε_F is the Fermi energy, and the remaining symbols have their usual meanings. The sum is over the two-dimensional electron wave vector $\mathbf{k}$. $S = 1$ corresponds to the unscreened case. The matrix elements are given by

$$| M(\mathbf{Q}) |^2 = \frac{Z^2 \hbar Q}{2\rho u_l} \tag{9}$$

for deformation-potential scattering,

$$| M(\mathbf{Q}) |^2 = \frac{\hbar (e h_{14})^2}{2\rho u_l} \frac{9 Q_\perp^2 Q_{||}^4}{2Q^7} \tag{10}$$

for longitudinal piezoelectric scattering, and

$$| M(\mathbf{Q}) |^2 = \frac{\hbar (e h_{14})^2}{2\rho u_t} \frac{8 Q_\perp^4 Q_{||}^2 + Q_{||}^6}{2Q^7} \tag{11}$$

for transverse piezoelectric scattering.

For the deformation-potential interaction, these equations can be reduced to

$$\begin{aligned} \left\langle \frac{\partial \varepsilon}{\partial t} \right\rangle = \frac{-2Z^2 (2m^*)^{1/2}}{\rho \varepsilon_F (2\pi)^2} & \int_0^\infty d\varepsilon_{\mathbf{k}} \frac{f(\varepsilon_{\mathbf{k}})}{\sqrt{4\varepsilon_{\mathbf{k}}}} \int_0^\infty dQ_{||} \int_{Q_{\perp\min}}^{Q_{\perp\max}} dQ_\perp \{ \exp[(1/kT_l - 1/kT_e)\hbar\omega_Q] - 1 \} \\ & \times N_Q I^2(Q_\perp) S^2(Q_{||}) (Q_\perp^2 + Q_{||}^2) \\ & \times \left[1 - \frac{(\hbar\omega_Q - \varepsilon_{Q_{||}})^2}{4\varepsilon_{\mathbf{k}} \varepsilon_{Q_{||}}} \right]^{-1/2} [1 - f(\varepsilon_{\mathbf{k}} + \hbar\omega_Q)] , \end{aligned} \tag{12}$$

where $Q_{\perp\min} = 0$ if $\varepsilon_{Q_{||}} < 4\varepsilon_k$ or if $(\varepsilon_{Q_{||}} - 2\sqrt{\varepsilon_k \varepsilon_Q})^2 < (\hbar u Q_{||})^2$, and

$$Q_{\perp\min} = [(\varepsilon_{Q_{||}} - 2\sqrt{\varepsilon_k \varepsilon_{Q_{||}}})^2 / \hbar^2 u_l^2 - Q_{||}^2]^{1/2} \tag{13}$$

otherwise. $Q_{\perp\max}$ is given by

$$Q_{\perp\max} = [(\varepsilon_{Q_{||}} + 2\sqrt{\varepsilon_k \varepsilon_{Q_{||}}})^2 / \hbar^2 u_l^2 - Q_{||}^2]^{1/2} . \tag{14}$$

The results for piezoelectric scattering are similar. Over certain temperature ranges, approximations can be made that lead to the simple temperature dependences discussed earlier.[10,16] For the temperatures used in this study, it was necessary to perform the integrations numerically for given T_e, T_l, and ε_F.

III. EXPERIMENT

The theory gives the energy-loss rate as a function of T_l, T_e, and Z. We measure T_e as a function of input power at known T_l and then adjust Z to produce the best fit between the theory and the experimental data. T_e is measured by using the well-known temperature dependence of the amplitudes of the Shubnikov–de Haas (SdH) oscillations.[19] The SdH oscillations occur when the magnetic field strength perpendicular to the plane of the 2D electron gas (2DEG) is varied, causing Landau levels to pass through the Fermi level, changing the density of final states available for scattering, and resulting in oscillations in the sample resistance. Increasing temperatures damp the oscillations through thermal broadening. This effect has frequently been used to measure T_e in bulk semiconductors and in Si inversion layers,[20,21] and was the

method used by Hirakawa and Sakaki in their study.

Although there is a theory to predict the amplitudes of the SdH oscillations as a function of temperature, it includes a number of unknown parameters, and is not accurate enough to be used directly. We therefore calibrate the temperature dependence of the SdH oscillations experimentally. This is done by measuring the amplitudes at a sequence of lattice temperatures while using low input powers to ensure that the electron heating is negligible. The electron temperature at higher input powers can then be determined by comparing the resulting amplitudes with the calibration curve.

The SdH amplitudes are an excellent measure of T_e because they have no direct dependence on T_l. A quantity which is a direct function of both T_l and T_e cannot be used because its T_e dependence cannot be calibrated. One recent study has used the T dependence of the sample resistance to determine T_e.[22] Resistance depends both on the concentration and on the mobility, both of which can depend differently on T_e and T_l, and this method must therefore be used with care. If, for example, the mobility is phonon limited, then it is virtually independent of T_e, and cannot be used to measure T_e. In most cases, the mobility will be a function of both T_e and T_l, and it will be difficult to calibrate the T_e dependence of the resistance. We do not believe resistance to be an accurate measure of T_e for our particular samples. We attempted to measure T_e using the sample resistance, but were unable to fit the resulting data with a consistent value of Z^2. Data taken at the higher input powers required a value of Z^2 that was an order of magnitude different from that required for the lower input powers. This is to be compared with the results obtained by using the SdH amplitudes as a measure of T_e, which required less than a 10% variation in the value of Z^2 needed to fit the data over the same range of input power.

A number of complications must be considered to ensure the accuracy of the results. The theoretical energy-loss rate with which we compare our experiment was calculated with the assumption that the density of states is constant. However, the SdH oscillations occur because the magnetic field has disturbed the density of states and hence the scattering rate. We must minimize the disturbance if the comparison is to be accurate. By using magnetic fields low enough that the total change in the sample resistance is small, we can be assured that the change in the scattering rate is minimal. The total resistance change for our samples over the range of magnetic field used (up to 4 T) is approximately 10%. We are able to use what would normally be considered a high magnetic field because we have used samples with relatively low mobility ($\sim 2\times 10^4$ cm^2/V sec), and correspondingly broad Landau levels.

It is also important to verify that the scattering theory used is applicable to the experimental conditions. For a degenerate semiconductor the use of the golden rule is not justified unless $\hbar/\varepsilon_F < \tau$, where τ is the time between collisions.[19] The Fermi energy for our samples in 16.5 meV, as determined from the period of the SdH oscillations, and we infer a τ of 8×10^{-13} from the mobility, which easily satisfies this inequality. Note also that the mean free path of an electron is longer than its de Broglie wavelength.

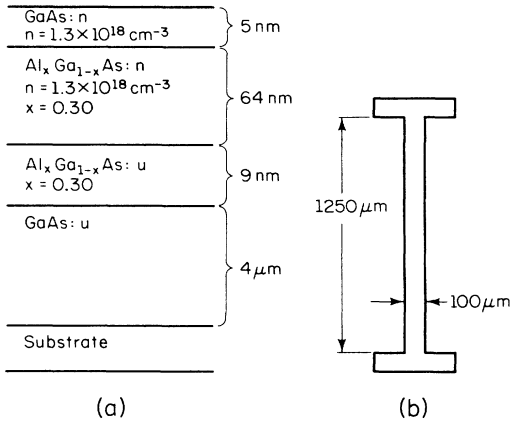

FIG. 1. (a) The structure of the sample used in this study. The sample was grown by atmospheric pressure MOCVD. (b) The mesa pattern for the devices measured.

The samples used in this study were grown on Cr-doped semi-insulating GaAs (100) substrates using atmospheric-pressure metal-organic chemical-vapor deposition (MOCVD). The structure used was a modulation doped single heterojunction, details of which are shown in Fig. 1(a). Photolithography was used to define mesas with a large aspect ratio, as shown in Fig. 1(b), and Ohmic contacts were made using AuGe/Ni/Au, alloyed at 430 °C for 30 sec in a flowing hydrogen atmosphere.

A dc current source was used to measure ρ_{xx}. The contact resistance is minimal in comparison with the sample resistance, which was approximately 2 $k\Omega/\square$. Lattice heating is expected to be unimportant at the input powers used. A simple calculation assuming that a thin slab of GaAs containing the conduction electrons is heated and that the heat diffuses to the substrate which remains at the temperature of the liquid He shows that the input powers used (less than 120 W/m^2) can support a temperature rise of less than 1 mK.

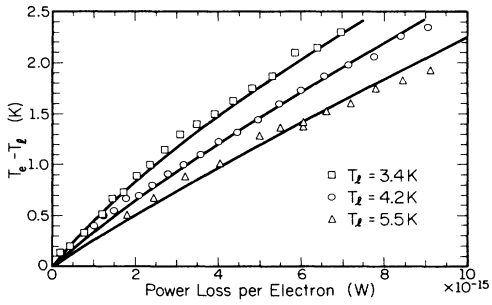

FIG. 2. The measured electron temperature vs the power loss at three different lattice temperatures. The solid lines are the power loss computed from the screened theory with a deformation potential of 16 eV, and include the loss to piezoelectric scattering.

Preliminary measurements were made on devices from across the sample wafer to ensure the consistency of the results, and detailed measurements, at a number of different lattice temperatures, were made on a limited number of devices. Figure 2 shows the dependence of $(T_e - T_l)$ on the power loss per electron for one device at various lattice temperatures.

IV. DISCUSSION

We determine the deformation potential by adjusting its value to produce the best fit between the experimentally determined energy-loss rate and the total (deformation potential plus piezoelectric) theoretical energy-loss rate. At the temperatures of our study, piezoelectric scattering contributes only about 10% of the total energy-loss rate, which minimizes any possible error arising from uncertainties in the value of the piezoelectric constant.

Using the self-consistent wave function, the screened theory requires a deformation-potential constant of $Z = 16.0 \pm 0.5$ eV, and the unscreened theory requires $Z = 11.5 \pm 0.5$ eV. Over the temperature range used in this study, the difference in the temperature dependence of the screened and unscreened theory was not great enough to allow us to reliably choose one or the other based on our data. This is contrary to what would be expected based on the analytic solution. The discrepancy arises because the approximations necessary to solve the integrals analytically are not applicable in the temperature range of this study (and they remain inaccurate to temperatures below 1 K). Hence, although the absolute value of the energy-loss rate is reduced by screening, the T dependence of the power loss is essentially unchanged, and either theory fits the data equally well. Figure 2 shows the best simultaneous fit of the energy-loss rate at three different lattice temperatures. Figure 3 shows how the fit is changed by a variation in Z. The use of the Fang-Howard variational wave function would reduce the theoretical energy-loss rates by approximately 30%, and required a corresponding increase in the value of Z^2 necessary to fit the experimental data.

We note that our measured energy-loss rate is almost twice that found by Hirakawa and Sakaki, and required a corresponding increase in Z^2. It is not clear why the energy-loss rates were different. If the phonon modes are disturbed by the proximity of the surface of the crystal, then the resulting energy-loss rates could depend sensitively on the distance between the electron sheet and the surface. It may also be significant that Hirakawa used samples grown by molecular beam epitaxy (MBE), while we used samples grown by MOCVD.

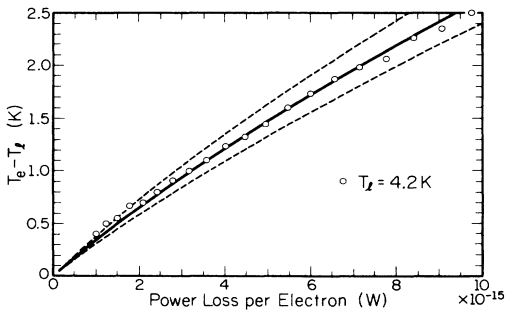

FIG. 3. This figure shows how a change in Z affects the accuracy of the fit. The middle line is the screened power loss chosen for the best fit (16 eV). The upper and lower lines are the power loss for $Z = 15$ and 17 eV, respectively.

We have observed that repeated thermal cycling of a sample between room temperature and liquid-helium temperatures results in an increase in the energy-loss rate. It is not clear what the origin of this increase is. One could speculate that thermal cycling may generate some defects which scatter carries inelastically. Excitation of electrons in interface states and subsequent phonon-assisted recombination is one possible process. For devices that had been cycled many times, the energy-loss rate for a given T_e was as much as 20% more than the original energy-loss rate. The data shown in Fig. 2 is from a device which had been cooled only once, over a period of several hours.

In contrast to previous work,[1,5–7] we do not interpret our data as proving that earlier measurements of the deformation-potential constant in bulk GaAs were in error. As Walukiewicz has pointed out,[4] the adoption of the larger values of the deformation-potential constant would limit the theoretical maximum of the 77 K mobility in bulk GaAs to less than the observed values. If we consider polar optical, piezoelectric, and deformation-potential scattering, but not ionized impurity scattering (which is sample dependent), then a deformation potential of 16 eV limits the mobility at 77 K to 1.37×10^5 cm^2/V sec. Even without the inclusion of ionized impurity scattering, which is significant, these values are considerably below the observed mobilities in high purity material, which have approached 2.10×10^5 cm^2/V sec.[5] Without ionized impurity scattering, a deformation potential of 11 eV would limit the mobility to 2.10×10^5 cm^2/V sec. One study of the deformation-potential constant in bulk GaAs, which is frequently cited to justify the assertion that the larger values of Z are correct, used bulk GaAs with impurity levels of $10^{16}/cm^3$ or more, and found $Z = 16$ eV.[23] At this impurity level, deformation-potential scattering is a small part of the total scattering, and a small inaccuracy in the estimate of the impurity scattering rates can result in a large error in the estimate of the deformation-potential scattering. The study by Wolfe, Stillman, and Lindley, which found $Z = 7$ eV, was done using material with an impurity level of less than 10^{14} cm^3, which reduces this problem.[8]

Studies of the mobility in silicon inversion layers provide an additional indication that the current theories of scattering in quasi-2D systems do not properly account for the deformation-potential interaction at interfaces. They have also found that the value of the deformation-potential constant necessary to fit the data was anomalously large.[24]

In summary, we have measured the energy-loss rate of electrons in a modulation-doped single heterojunction at several lattice temperatures. The theoretical energy-loss rates are found to depend sensitively on the wave function perpendicular to the interface. Using a self-consistent

wave function, we find that the data can be fit by either a screened theory, which requires $Z = 16.0$ eV, or by an unscreened theory, which requires $Z = 11.5$ eV. We suggest that these values not be regarded as measures of the deformation potential in bulk GaAs. Discrepancies between the bulk value for Z and the value measured in heterojunctions may stem from additional scattering mechanisms in heterojunctions which have not been accounted for, or from the differences in the phonon modes, caused either by the $Al_xGa_{1-x}As/GaAs$ interface, or by the proximity of the surface.

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