

the Elliott theory for indirect transitions² according to the energy-band structure for the $\text{GaAs}_x\text{P}_{1-x}$ mixed-crystal system given in Ref. 1 and the GaP absorption data in Ref. 4, is $\sim 1.5 \text{ cm}^{-1}$. This calculated value is in good agreement with the measured absorption excess in the LA_B phonon-assisted component shown in Fig. 4.

IV. SUMMARY

Pronounced optical absorption appears just above the exciton energy gap E_{gx} of the indirect-gap semiconductor GaP in the presence of a few mole % of GaAs. The most prominent absorption component is due to the creation of free excitons and free-electron-hole pairs. The threshold for the latter process is not clearly defined in the mixed crystals. Momentum is conserved in this transition through scattering in the localized potential of the arsenic impurities. The decrease in the energy difference between the Γ_1 and X_1 conduction-

band minima in $\text{GaAs}_x\text{P}_{1-x}$ compared with GaP also induces extra indirect absorption in the mixed crystal associated with the emission or absorption of momentum-conserving LA phonons. The free-particle absorption, measured at an arbitrarily chosen energy 0.04 eV above E_{gx} , is ~ 7 times larger for each arsenic atom than for the phosphorus atom it replaces, but is only $\sim 0.1\%$ of the corresponding absorption due to nitrogen substituents in GaP. These results demonstrate the high sensitivity of the optical absorption process to the difference between the electronic properties of the host and substituent atoms.

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Note on Optical Mixing by Mobile Carriers in Semiconductors*

B. S. KRISHNAMURTHY AND V. V. PARANJAPE

Department of Physics, Lakehead University, Port Arthur, Ontario, Canada

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It is shown that nonlinearity due to energy dependence of the carrier momentum relaxation time in semiconductors may arise even when the electrons have the same temperature as the lattice. This leads to modifications in the results of a recent paper.

IN a recent paper by Kaw,¹ it has been shown that the energy dependence of the carrier momentum relaxation time may contribute to optical mixing even in parabolic-band-type semiconductors. In this paper we point out that the model used by the author¹ is valid when electron temperature T_e is very much greater than lattice temperature T , while in the warm-electron situation $(T_e - T)/T \ll 1$ (as in Ref. 1) the calculations require modifications.

We start with the equation of motion for an average carrier in the form

$$\partial \mathbf{V} / \partial t + \mathbf{V} / \tau(\bar{\epsilon}) = -e\mathbf{E} / m, \quad (1)$$

where $\mathbf{E} = \mathbf{E}_1 \exp i\omega_1 t + \mathbf{E}_2 \exp i\omega_2 t + (\text{complex conjugate})$ is the electric field due to laser beams, \mathbf{V} is the average carrier velocity, m is the effective mass, and $\tau(\bar{\epsilon})$ is the carrier momentum relaxation time which depends on the average energy of the carrier $\bar{\epsilon}$. In the presence of laser beams, the average energy at time t may be written

$$\bar{\epsilon} = \frac{3}{2} k T_e + \frac{1}{2} m V^2, \quad (2)$$

where the second term in the above equation is the energy associated with the drift velocities of the carriers. This term is usually disregarded if electrons are hot, i.e., $T_e \gg T$, but in the case of warm electrons, it has been demonstrated by Paranjape and Alba² that the term is important and can lead to electron cooling by a dc electric field. Here it is shown that inclusion of the term produces significant changes in the results of Ref. 1. In the notations of Ref. 1 we write

$$1/\tau = (1/\tau_0) \{ 1 - \frac{1}{2} n [(T_e/T - 1) - mV^2/3kT] \}, \quad (3)$$

where n is an integer.

The electron temperature T_e is determined by the energy balance equation which in view of Eq. (2) is given by

$$d\bar{\epsilon}/dt = -e\mathbf{V} \cdot \mathbf{E}, \quad (4)$$

where we have neglected the energy loss of the electrons to the lattice. This simplifying assumption as in Ref. 1 is correct provided ω_1 , ω_2 , and $\omega_1 - \omega_2$ are very much greater than the reciprocal of the energy relaxation time. We now develop \mathbf{V} and T as a sum of com-

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¹ P. Kaw, Phys. Rev. Letters 21, 539 (1968).

² V. V. Paranjape and E. de Alba, Proc. Phys. Soc. (London) 85, 945 (1965).

ponents of frequencies ω_1 , ω_2 , $2\omega_1$, etc. From Eqs. (1) and (3), we obtain the various velocity components

$$V(\omega_1) = -eE_1\tau_0/m(1+i\omega_1\tau_0), \quad (5)$$

$$V(\omega_2) = -eE_2\tau_0/m(1+i\omega_2\tau_0), \quad (6)$$

and

$$V(2\omega_1-\omega_2) = (n/2\tau_0^2kT)[i(2\omega_1-\omega_2)\tau_0+1]^{-1} \times [mV^2(\omega_1)V^*(\omega_2)+V(\omega_1)kT_e(\omega_1-\omega_2)+V^*(\omega_2)kT_e(2\omega_1)]. \quad (7)$$

If one neglects in Eq. (7), the changes in the electron temperature, i.e., $T_e=T$, we get for the output of the current component with frequency $2\omega_1-\omega_2$ using Eqs. (5) and (6) in the following form:

$$J_0(2\omega_1-\omega_2) = -NeV(2\omega_1-\omega_2) = -\frac{1}{2}n[Ne^4E_1^2E_2/m^2kT\omega_1\omega_2(2\omega_1-\omega_2)(\omega_1\tau_0)], \quad (8)$$

where $\omega_1\tau_0 \gg 1$ and N is the carrier concentration. In Ref. 1, for $T_e=T$ there is no contribution to $J(2\omega_1-\omega_2)$. We shall now establish that the first term in Eq. (7) which leads to Eq. (8) is larger in its contribution to the current component $J(2\omega_1-\omega_2)$ than the other two terms. Using Eqs. (1)-(4), the temperature components are expressed by

$$T_e(2\omega_1) = -[e^2E_1^2\tau_0/3(1+\omega_1^2\tau_0^2)^2\omega_1mk] \times [2\omega_1\tau_0+i(1-\omega_1^2\tau_0^2)] \quad (9)$$

and

$$T_e(\omega_1-\omega_2) = -4e^2E_1E_2^*/3km(1+\omega_1^2\tau_0^2)(1+\omega_2^2\tau_0^2)(\omega_1-\omega_2) \times [\tau_0(\omega_1-\omega_2)-i(1+\omega_1\omega_2\tau_0^2)]. \quad (10)$$

It should be noted that Eqs. (9) and (10) differ from Eqs. (7a) and (7b) of Ref. 1, due to the inclusion of the drift term in Eq. (2). Combining Eqs. (5)-(7), (9), and (10), we write

$$J(2\omega_1-\omega_2) = \frac{1}{2}n \frac{Ne^4E_1^2E_2}{m^2kT\omega_1\omega_2(2\omega_1-\omega_2)} \times \left[-\frac{1}{\omega_1\tau_0} + i \frac{A}{(\omega_1\tau_0)^2} \right] = J_0(2\omega_1-\omega_2)[1-i(A/\omega_1\tau_0)], \quad (11)$$

where

$$A = \frac{6\omega_1^3-5\omega_2^3-30\omega_1^2\omega_2+25\omega_1\omega_2^2}{3(\omega_1-\omega_2)(2\omega_1-\omega_2\omega_2)}.$$

We have assumed $\omega\tau_0 \gg 1$ and retained terms up to $1/(\omega\tau_0)^2$. The first and the second term in the square bracket of Eq. (11) arise due to the electron drift and the changes in the electron temperature, respectively. It therefore follows that the contribution to the non-linearity due to the drift term may be important when $(A/\omega_1\tau_0) \ll 1$.

Optical Spectra of Donor Impurities in InSb in High Magnetic Fields

R. KAPLAN

Naval Research Laboratory, Washington, D. C. 20390

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Experimental studies of the donor-impurity excitation spectra in InSb in high magnetic fields have yielded good qualitative agreement with theoretical predictions with regard to transition energies, selection rules, polarization, and absorption constants. Quantitatively, the data demonstrate the need for including the effects of nonparabolicity in the theory, and indicate that corrections to the calculated ground-state energy are required. The usefulness of optical experiments in studying carrier freeze-out effects is briefly discussed.

I. INTRODUCTION

THE effects of a magnetic field on the energy levels and excitation spectra of impurity atoms in semiconductors are usually considered in the limiting cases of very small and very large magnetic fields. In the former case the effects appear as small shifts and splittings of the zero-field levels and optical transitions, while in the latter case there is a drastic qualitative change in the level structure and the excitation spectra.

A number of theoretical calculations¹⁻⁴ for the high-field case have succeeded in giving a detailed account of the impurity states and optical transitions for donor

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⁴ H. Hasegawa, in *Physics of Solids in Intense Magnetic Fields*, edited by E. D. Haidemenakis (Plenum Press, Inc., New York, 1969).