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Parametric Amplification of Ultrasound in Strong Magnetic Fields

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When ultrasound propagates in a piezoelectric semiconductor, the acoustoelectric interaction between the acoustic flux and the conduction electrons leads to parametric amplification of acoustic waves of frequencies different from that of the initial ultrasonic wave. The acoustoelectric interaction between the electrons and the ultrasound can be modified by applying a strong magnetic field transverse to the direction of propagation of the ultrasound. The application of the magnetic field is shown to lead to the enhancement of the parametric-amplification process for waves whose frequency lies in the vicinity or below the frequency of maximum linear gain (or loss) in high-mobility semiconductors such as InSb or GaAs. However, the threshold strain necessary to have the parametric amplification exceed the linear loss also increases with magnetic field owing to the large enhancement of the linear electronic losses with magnetic field.

I. INTRODUCTION

In a recent paper Conwell and Ganguly[1] have discussed mixing of acoustic waves in piezoelectric semiconductors due to the acoustoelectric interaction between the waves and the conduction electrons. The frequency mixing occurs because of the interactions between electrons which are bunched by the piezoelectric fields that accompany the acoustic flux. Therefore, any process which enhances the bunching of the carriers by the piezoelectric field will enhance the parametric amplification of acoustic flux while any process which decreases the bunching will also decrease the nonlinear gain due to parametric amplification. How well the electrons are bunched by the piezoelectric field depends upon how fast they can follow the field, i.e., their mobility, and how fast the density gradients resulting from the bunching are smoothed out as a result of diffusion.

If a strong magnetic field is applied transverse to the direction of propagation of the ultrasound, both the mobility and the diffusion coefficient of the carriers are reduced. This reduction is particularly evident in high-mobility semiconductors such as InSb and GaAs. Whether the application of the magnetic field will enhance or decrease bunching will depend upon whether the reduction of the mobility, which prevents the carriers from following the piezoelectric field and therefore leads to debunching, is more important than the reduction of diffusion which enhances the bunching. If it is more important, then the presence of the field will decrease the parametric amplification of acoustic waves, while, if the reduction of diffusion is more important, the parametric processes will be enhanced.

The effect of a transverse magnetic field on the linear gain or loss due to the acoustoelectric interaction has been predicted[2] and observed.[3,4] In a previous paper, the author[5] calculated the effect of the transverse magnetic field on second-harmonic generation and found that the field modified both the magnitude and the frequency dependence of the harmonic generation. In this paper, we present a more detailed theory of the mixing of acoustic waves in a transverse magnetic field following essentially along the same lines as Ref. 1 and a previous paper.[5] We then apply our results to a discussion of subharmonic generation in high-mobility piezoelectric semiconductors such as InSb and GaAs.

II. THEORY

The acoustoelectric interaction between ultrasound and the conduction electrons can be described by the equation of motion of the lattice, the piezoelectric equations of state, Maxwell's equations, and an equation for the electronic current induced by the ultrasound. Since acoustic waves which induce transverse piezoelectric fields have a much weaker coupling to the electrons than those which induce longitudinal piezoelectric fields,[5] we need only Poisson's equation and the equation of continuity to determine the electric displacement in terms of the electronic current induced by the wave. The relevant set of equations is

$$\rho \frac{\partial^2 \xi_i}{\partial t^2} = \frac{\partial T_{ij}}{\partial X_j} , \tag{1}$$

$$T_{ij} = C_{ijkl} S_{kl} - \beta_{ijk} E_k , \tag{2}$$

$$D_i = \epsilon E_i + 4\pi \beta_{ijk} S_{jk} , \tag{3}$$

$$S_{ij} = \frac{1}{2}\left(\frac{\partial \xi_j}{\partial X_i} + \frac{\partial \xi_i}{\partial X_j}\right), \tag{4}$$

$$\vec{\nabla} \cdot \vec{D} = -4\pi e (n - n_0) , \tag{5}$$

$$-e\frac{\partial n}{\partial t}+\vec{\nabla}\cdot\vec{J}=0\ . \tag{6}$$

This set of equations needs to be supplemented by an equation for the electronic current $\vec{J}$. Under conditions where $\omega_0\tau \ll 1$ and $ql \ll 1$ or $\omega_0\tau \gg 1$ and $qR \ll 1$, where ω_0 is the cyclotron frequency, τ the electron relaxation time, l the electron mean free path, and R the Larmor radius, we can write down the following phenomenological equation for the current density:

$$\vec{J}=ne\mu\vec{E}+(\mu/c)\vec{J}\times\vec{B}-eD\vec{\Delta}n. \tag{7}$$

Following along the same lines as Ref. 1, we write

$$\vec{J}=\vec{J}_0+\textstyle\sum_m \vec{J}_m\, e^{i(\vec{k}_m\cdot\vec{r}-\omega_m t)}\ , \tag{8a}$$

$$\vec{E}=\vec{E}_0+\textstyle\sum_m \vec{E}_m\, e^{i(\vec{k}_m\cdot\vec{r}-\omega_m t)}\ , \tag{8b}$$

with similar equations for the carrier density, magnetic field, and strain. Using (8) in Eqs. (1)–(7) and neglecting the transverse electric fields, we obtain the following relation which determines J_{mx}:

$$J_{mx}=\frac{\sigma_0E_{mx}+(\mu/V_s)\sum_{m'}(J_{m'x}E_{m-m'x}e^{i\vec{\Delta}_{m-m'}\cdot\vec{r}}+J_{m'+mx}E_{m'x}e^{i\vec{\Delta}_{m+m'}\cdot\vec{r}})}{\gamma b+\omega_m/\omega_D}\ . \tag{9}$$

Here

$$\vec{\Delta}_{m\pm m'}=\vec{k}_{m\pm m'}\mp\vec{k}_{m'}-\vec{k}_m\ ,\quad b=1+(\omega_0\tau)^2\ ,$$

$$\gamma=1-V_d/V_s\ ,$$

σ_0 is the dc conductivity, ω_D is the diffusion frequency, and V_d is the electron drift velocity. Equation (9) is valid for both the Hall and Corbine geometries except that in the latter case we have the extra condition[5] that $\omega_0\tau \gg 1$ and of course the form of V_d differs for the two situations.

We now wish to limit ourselves to the case where we have three waves present where the highest frequency $\omega_3=\omega_1+\omega_2$. Since we are interested in the case of weak interactions, we assume that the relation between the fields and the strains at a frequency ω_m is the same as in the absence of the interaction. We shall also restrict ourselves to treating the situation in which the three waves are collinear. It has been shown that in the presence of linear gain or loss, the importance of phase matching decreases. Writing the piezoelectric field as the first-order field plus a correction due to the interaction between waves, we have

$$E_{mx}=E^1_{mx}+E^2_{mx}\ , \tag{10a}$$

$$E^1_{mx}=-\frac{4\pi\beta}{\epsilon}\left(\frac{\gamma b+i\omega_m/\omega_D}{\Gamma_m}\right)S_m\ , \tag{10b}$$

$$E^2_1=-\frac{i\mu}{V_s}\left(\frac{4\pi\beta}{\epsilon}\right)^2\frac{\omega_c}{\omega_1}\left(\frac{2\gamma b+i\omega_1/\omega_D}{\Gamma_1\Gamma^*_2\Gamma_3}\right)S^*_2S_3e^{i\Delta x}\ , \tag{10c}$$

$$E^2_2=-\frac{i\mu}{V_s}\left(\frac{4\pi\beta}{\epsilon}\right)^2\frac{\omega_c}{\omega_2}\left(\frac{2\gamma b+i\omega_2/\omega_D}{\Gamma^*_1\Gamma_2\Gamma_3}\right)S^*_1S_3e^{i\Delta x}\ , \tag{10d}$$

$$E^2_3=-\frac{i\mu}{V_s}\left(\frac{4\pi\beta}{\epsilon}\right)^2\frac{\omega_c}{\omega_3}\left(\frac{2\gamma b+i\omega_3/\omega_D}{\Gamma_1\Gamma_2\Gamma_3}\right)S_1S_2e^{-i\Delta x}\ . \tag{10e}$$

Here

$$\Gamma_m=\gamma b+i(\omega_m/\omega_D+\omega_c/\omega_m)\ ,\quad \Delta=k_3-k_2-k_1\ ,$$

β is the relevant component of the piezoelectric tensor, and $\omega_c=4\pi\sigma_0/\epsilon$ is the dielectric relaxation frequency. Except for factors of 2 which arise from a difference in the definition of our expansion coefficients in (8), our results go over to those previous obtained in the limit of zero magnetic field ($b=1$).

Following the same procedure as in Ref. 1, we write the sound-wave amplitude in the form

$$\xi_i=U_i(x)e^{i(k_ix-\omega_it)}\ . \tag{11}$$

Because of the smallness of the electromechanical coupling constant, the amplitude will change very little over the distance of a wavelength, and we obtain the following set of first-order equations for U_1, U_2, and U_3:

$$\frac{dU_1}{dx}=-\alpha_1U_1-\eta_1U^*_2U_3e^{i\Delta x}\ , \tag{12a}$$

$$\frac{dU_2}{dx}=-\alpha_2U_2-\eta_2U^*_1U_3e^{i\Delta x}\ , \tag{12b}$$

$$\frac{dU_3}{dx}=-\alpha_3U_3+\eta_3U_1U_2e^{-i\Delta x}\ , \tag{12c}$$

where the absorption coefficient α_i has been derived previously[5] and

$$\eta_1=\frac{i\mu}{2V_s}k_2k_3\frac{\omega_c}{\omega_1}\left(\frac{4\pi\beta}{\epsilon}\right)^2\frac{\beta}{C}\left(\frac{2\gamma b+i\omega_1/\omega_D}{\Gamma_1\Gamma^*_2\Gamma_3}\right), \tag{13a}$$

$$\eta_3=\frac{i\mu}{2V_s}k_1k_2\frac{\omega_c}{\omega_3}\left(\frac{4\pi\beta}{\epsilon}\right)^2\frac{\beta}{C}\left(\frac{2\gamma b+i\omega_3/\omega_D}{\Gamma_1\Gamma_2\Gamma_3}\right)\ . \tag{13b}$$

The expression for η_2 is identical to that for η_1 with the subscripts 1 and 2 interchanged. Using the set of coupled equations in Eq. (13), we are now in a position to discuss various parametric processes.

A. Down Conversion

When the amplitude of the acoustic wave at the highest of the three frequencies ω_3 is large while the amplitude at both ω_1 and ω_2 is small, we can decouple Eq. (12c) by neglecting the second term on the right-hand side as being small compared to the first term. This neglects the depletion of the pump and is therefore only valid at small x. As a result, the second terms on the right-hand side of Eqs. (12a) and (12b) act as sources at the frequencies ω_1 and ω_2 thereby leading to nonlinear gain at the lower frequencies. Making these approximations, the solution to Eq. (12c) is

$$U_3(x) = U_3(0)\, e^{-\alpha_3 x} ,$$

while the solutions to Eq. (12a) and (12b) are

$$U_1(x) = (1/2m)\exp[-\tfrac{1}{2}(\alpha_1+\alpha_2)x]e^{i\Delta x/2}(\{\eta_1 U_2^*(0)U_3(0) + [\tfrac{1}{2}(\alpha_2-\alpha_1-i\Delta)]U_1(0)\}\sinh mx + 2mU_1(0)\cosh mx) , \qquad (14a)$$

$$U_2^*(x) = (1/2m)\exp[-\tfrac{1}{2}(\alpha_1+\alpha_2)x]e^{-i\Delta x/2} \times(\{\tfrac{1}{2}(\alpha_2-\alpha_1-i\Delta)U_2^*(0) - \eta_2^* U_3^*(0)U_1(0)\}\sinh mx - mU_2^*(0)\cosh mx) , \qquad (14b)$$

where

$$m = \{[\tfrac{1}{2}(\alpha_2-\alpha_1-i\Delta)]^2 + \eta_1\eta_2^*|U_3(0)|^2\}^{1/2}.$$

Equation (14) has been derived under the assumption the $\alpha_3 x \ll 1$. When this assumption is relaxed, the solution becomes

$$U_1(x) = \exp[-\tfrac{1}{2}(\alpha_1+\alpha_2+\alpha_3)x] \times\left[AI_{(1/2)[1+(\alpha_2-\alpha_1)/\alpha_3]}\left([\eta_1\eta_2^*]^{1/2}\left|\frac{U_3(0)}{\alpha_3}\right| e^{-\alpha_3 x}\right) + BI_{-1/2[1+(\alpha_2-\alpha_1)/\alpha_3]}\left([\eta_1\eta_2^*]^{1/2}\left|\frac{U_3(0)}{\alpha_3}\right| e^{-\alpha_3 x}\right)\right]. \qquad (15)$$

where $I_n(x)$ is the imaginary Bessel function of order n and argument x.[6] In the degenerate case where $\omega_1=\omega_2$ and we have subharmonic generation, the latter result simplifies to

$$U_1(x) = \tfrac{1}{2}e^{-\alpha_1 x}\left[\left(U_1(0) - \frac{\eta S_3(0)}{|\eta|\,|S_3(0)|}\,U_1{}^*(0)\right) \times e^{|\eta||S_3(0)|F(x)} + \left(U_1(0) + \frac{\eta S_3(0)}{|\eta|\,|S_3(0)|}\,U_1{}^*(0)\right) \times e^{-|\eta||S_3(0)|F(x)}\right], \qquad (16)$$

where

$$F(x) = (e^{-\alpha_3 x}-1)|\alpha_3| \quad \text{and} \quad \eta = \eta_1/k_3 .$$

The behavior of the above solutions has been discussed in great detail in Ref. 1. It is found that if the real part of m is larger than the imaginary part, then for large enough x, but with $\alpha_3 x \ll 1$,

$$U_1(x) \simeq C\, e^{\alpha_{\mathrm{NL}} x} , \qquad (17)$$

where the nonlinear gain coefficient is

$$\alpha_{\mathrm{NL}} = -\tfrac{1}{2}(\alpha_1+\alpha_2) + \mathrm{Re}\, m .$$

In particular, for the degenerate case $\omega_3=\omega=2\omega_1=2\omega_2$,

$$\alpha_{\mathrm{NL}} = -\alpha_1 + |\eta|\,|S_3(0)| , \qquad (18)$$

where

$$|\eta| = \frac{\mu}{2V_s{}^2}\,\frac{(\omega_c/c)(4\pi\beta/\epsilon)^2\beta}{(\gamma b)^2 + (\omega/2\omega_D + 2\omega_c/\omega)^2} \times\left(\frac{4(\gamma b)^2 + (\omega/2\omega_D)^2}{(\gamma b)^2 + (\omega/\omega_D+\omega_c/\omega)^2}\right)^{1/2} . \qquad (19)$$

If the drift velocity V_d is large enough so that $\alpha_1 < 0$, then Eq. (18) yields both a linear and a nonlinear gain for the subharmonic $\omega_1=\tfrac{1}{2}\omega$. On the other hand, if $\alpha_1>0$, then there will only be a net gain if $|\eta|\,|S_3(0)| > \alpha_1$. Thus for small strains at the pump frequency ω, there will be no net gain while at larger strains there will be, depending upon the value of $|\eta|$ for the particular values of the parameters involved.

In the absence of the magnetic field, it was found that the rate of the nonlinear gain was greatest at the subharmonic for a wide range of pump frequencies ω. We expect this to also be the case in the presence of magnetic fields since the subharmonic was favored to higher pump frequencies when $\gamma \neq 0$ than is the case when $\gamma = 0$.[1] As can be seen from Eqs. (12) and (19), the presence of the magnetic field can be taken into account in the expressions for η_i by replacing γ by γb. Therefore, the effect of the magnetic field is formally the same as having a larger value of γ. Of course, it would require much larger electric fields to get a γb value which could be obtained by applying only moderate magnetic fields in a high-mobility semiconductor.

Since the gain is expected to be greatest for subharmonic generation, we will limit ourselves to a discussion of the dependence of $|\eta|$ on magnetic field and pump frequency. In high-mobility semiconductors such as n-type InSb or GaAs, the dielectric relaxation frequency ω_c is much greater than the diffusion frequency ω_D. Under these conditions $|\eta|$ has a maximum as a function of magnetic field. This can be seen from Fig. 1 in which the ratio $|\eta(B)|/|\eta(0)|$ is plotted as a function of magnetic field for various frequencies. The parameters used in Fig. 1 are characteristic of n-type InSb at 77 °K. For frequencies in the vicinity or below the frequency of maximum linear gain, $\omega_m = (\omega_c\omega_D)^{1/2}$, $|\eta|$ is enhanced by factors of up to several orders of magnitude by the application of the magnetic field. In Fig. 2, $|\eta|$ is shown as a function of pump frequency for various values of b.

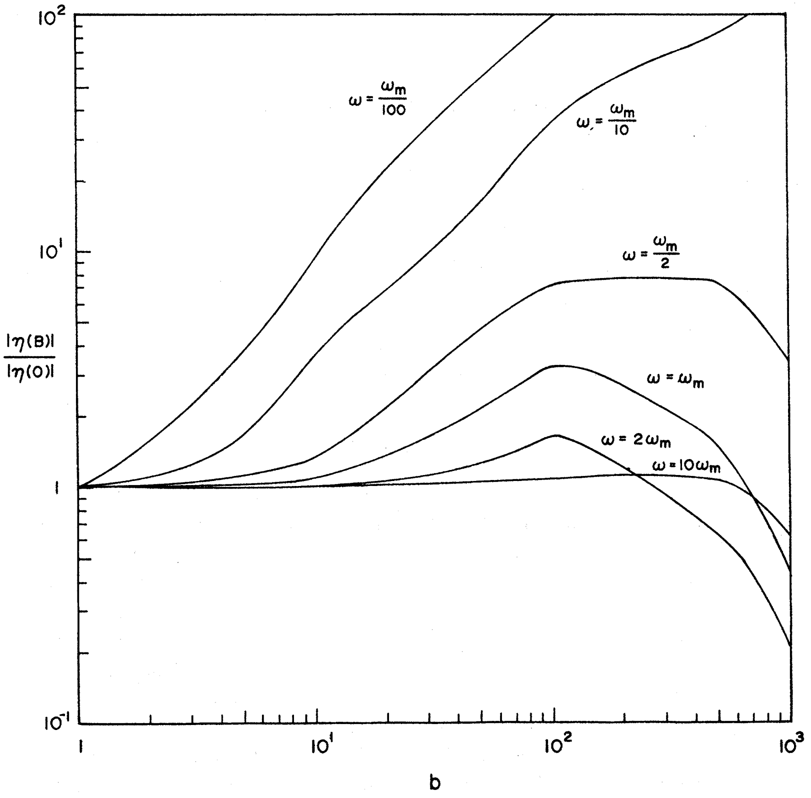

FIG. 1. Ratio $|\eta(B)|/|\eta(0)|$ is shown as a function of magnetic field for various frequencies and with $|\gamma| = 1$. The values of $\omega_c = 2.5 \times 10^{12}$ rad/sec and $\omega_D = 2.5 \times 10^8$ rad/sec characteristic of n-type InSb at 77 °K with a mobility of 10^5 cm^2/V sec and a carrier density of 2.5×10^{14} cm^{-3} have been used in making the calculation.

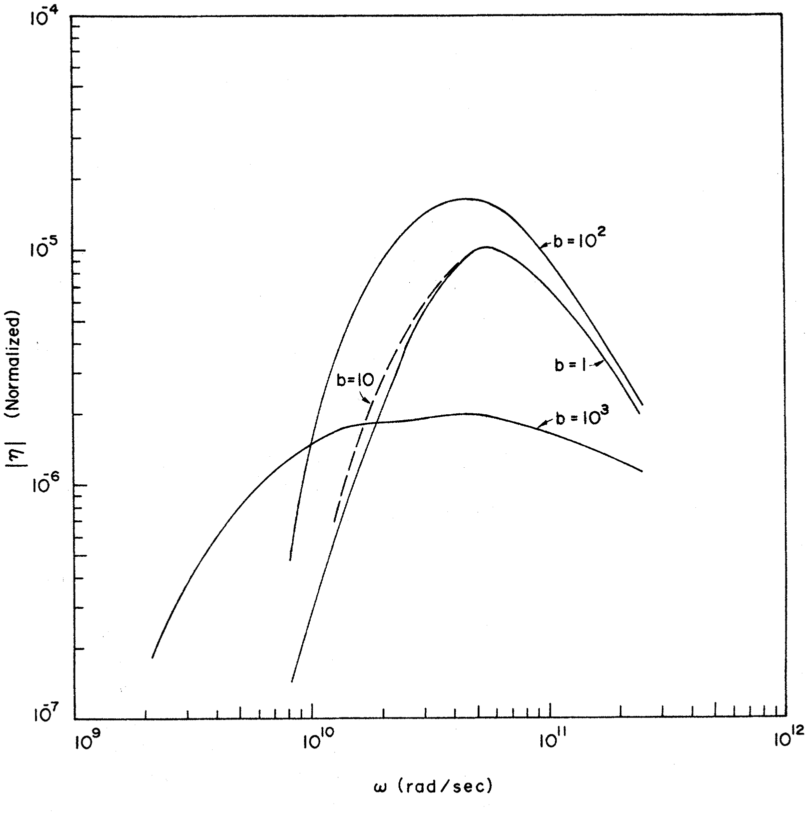

FIG. 2. Normalized nonlinear gain $|\eta| [(\mu/2\rho V_s^4)\omega_c(4\pi\beta/\epsilon)^2\beta]^{-1}$ is shown as a function of frequency for various values of b and with $|\gamma| = 1$. For n-type InSb, the normalization coefficient is 8×10^8.

In the absence of the magnetic field, $|\eta|$ has a maximum at a pump frequency about a little over twice the frequency of maximum linear gain. As the magnetic field initially increases, the curves are shifted upwards with the increase being greater on the low-frequency side of the peak than on the high-frequency side. This is due to the increased bunching which occurs as the field reduces the diffusion of the electrons out of the minima in the piezoelectric potential. Finally, at higher magnetic fields, the curve is shifted downwards and the peak is flattened out as the electrons are no longer able to follow the piezoelectric fields due to the reduction of their mobility.

For n-type InSb with a mobility at 77 °K of 10^5 $\text{cm}^2/\text{V sec}$, $b = 10^2$ corresponds to a magnetic field of 10 kG, while for n-type GaAs with a lower mobility of 8000 $\text{cm}^2/\text{V sec}$, $b = 10^2$ corresponds to a magnetic field of 125 kG. With even lower-mobility semiconductors such as CdS with a mobility of 300 $\text{cm}^2/\text{V sec}$, any significant enhancement of subharmonic generation would only occur at magnetic fields in the megagauss region. Since over a wide range of frequencies, the maximum enhancement of the subharmonic generation occurs from $b = 10^2$ to 10^3, one would only expect significant effects in high-mobility semiconductors where the peak in the subharmonic generation with field would occur at reasonable field strengths.

Finally, in Fig. 3 we show the magnitude of the threshold strain at the pump frequency, which is needed to have the nonlinear gain at the subharmonic exceed the linear electronic losses for $\gamma = 1$ ($V_d = 0$) as a function of b. Although the nonlinear gain increases with magnetic field, the linear electronic losses increase even more rapidly with magnetic field as shown in Fig. 4. As a result, the threshold strain for net amplification increases with magnetic field.

B. Upconversion

If initially there are waves present at frequencies ω_1 and ω_2, then a third wave will be generated at the frequency ω_3. At small enough x where the amplitude of the wave of frequency ω_3 is still small, we can neglect the terms containing U_3 in (12a) and (12b) to solve for U_1 and U_2;

$$U_1(x) = U_1(0)\, e^{-\alpha_1 x}\,, \quad U_2(x) = U_2(0)\, e^{-\alpha_2 x}\,. \tag{20}$$

Using this solution in Eq. (12c), we find that

$$U_3(x) = \frac{\eta_3 U_1(0) U_2(0)}{\alpha_3 - \alpha_1 - \alpha_2 - i\Delta} \times \{ \exp[-(\alpha_1 + \alpha_2 + i\Delta)x] - e^{-\alpha_3 x} \}\,. \tag{21}$$

When $\omega_1 = \omega_2$, Eq. (21) reduces to the result we obtained for the case of second harmonic generation.[5] Since this case has been discussed previously, we will not discuss it further in this paper.

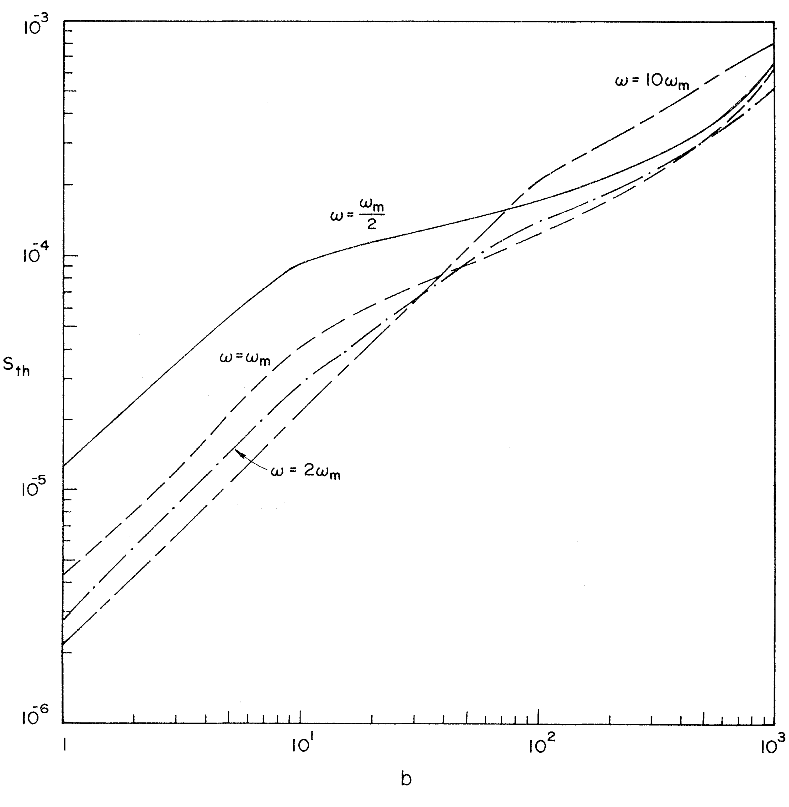

FIG. 3. Threshold strain for the net amplification of the subharmonic vs b for $\gamma = 1$ and the parameters characteristic of n-type InSb is shown for various pump frequencies.

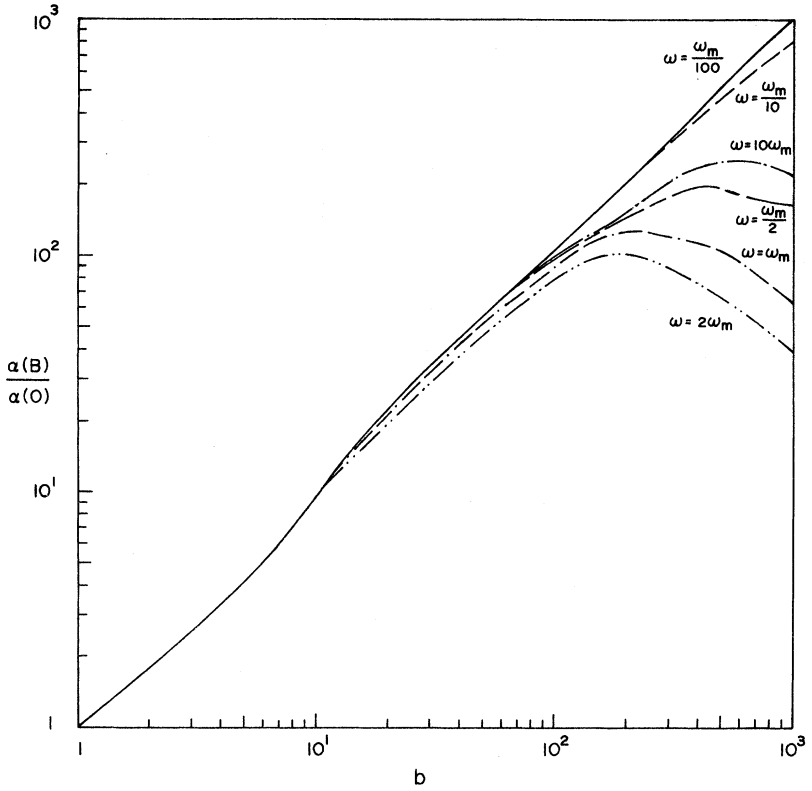

FIG. 4. Ratio $\alpha(B)/\alpha(0)$ at the subharmonic is shown as a function of magnetic field for various pump frequencies and for $\gamma = 1$.

III. DISCUSSION

In Sec. II, we have shown that the application of a dc magnetic field will lead to an enhancement of the nonlinear gain at the subharmonic frequency. From the general form of Eqs. (13) and (14), we expect a similar enhancement of all parametric processes with magnetic field. However, as can readily be seen by comparing Figs. 1 and 4, the enhancement of the linear electronic losses by the magnetic field is much greater than the enhancement of the nonlinear gain. Therefore, one expects that the effect of subharmonic generation to become important at lower pump strains in the absence of the field than in its presence. However, the enhancement of parametric effects with magnetic field would still occur when the frequency being parametrically amplified was not initially present such as in the case of second-harmonic generation.[5]

The growth of acoustic flux in a magnetic field has recently been observed via Brillouin scattering in GaAs at room temperatures.[7] In the weak-flux regime, the measurements were in excellent agreement with the predictions of the small signal theory. However, at higher-flux level, it was found that the enhancement of the growth rate with field decreased below that predicted by the small signal theory and finally in the strong-flux regime, the flux was found to saturate at intensities which were essentially independent of magnetic field. The theory presented here might account for some of the features in the intermediate-flux regime. Because of the large linear gain in the magnetic field, the flux will grow more rapidly than in the absence of the field. At flux levels where the nonlinear gain became more important than the linear electronic gain, the growth rate would not be as strongly dependent on the field as in the weak-flux regime since $|\eta|$ has a weaker magnetic field dependence than α, especially in the vicinity of the frequency ω_m. Finally, when the strong-flux regime is attained one would not expect the present theory to be valid since it assumes the linear growth rate is unaffected by the nonlinear interactions. Alternatively, since the linear growth rate is so enhanced by the field, the flux may grow into the strong-flux regime before pump strains at which the parametric-amplification effects become important are attained.

Note added in proof. In a private communication, Professor Bray has pointed out that in the experiments reported in Ref. 7, the value of γb in the presence of the magnetic field was chosen to be equal to the value of γ in the absence of the field. Since in our theory, the magnetic field comes in only through the parameter b, the enhancement of the growth rate should have been the same both with an without the magnetic field. The fact that the rates differed in the experiment indicates that in the strong flux regime, the theory presented here must be modified.

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Absorption from Neutral Acceptors in GaAs and GaP

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We present a new calculation of the absorption due to transitions of holes between neutral acceptors and the various valence-band sublevels in GaAs and GaP. The acceptor wave function was approximated by a previously suggested expression for ground-state wave functions appropriate to complicated band extrema. Numerical calculations of the absorption from inter-valence-band transitions of free holes and neutral acceptors have been performed. Good agreement with experimental results is obtained.

I. INTRODUCTION

Optical transitions of free holes between pairs of the three sublevels of the valence band in GaAs and GaP give rise to a characteristic infrared absorption.[1-3] In GaAs at room temperature, the absorption has qualitatively the same shape as the absorption due to the same transitions among the valence-band levels of Ge.[2,4-7] Balslev[2] calculated the room-temperature absorption in *p*-type GaAs and obtained good agreement with experimental results. No interpretation of the absorption at 77 °K[1] has been published. Recently, Wiley and DiDomenico (WD)[3] measured the infrared absorption of *p*-type GaP. In contrast to GaAs, the absorption spectrum of GaP has no significant structure. WD made a calculation of the room-temperature absorption, based on the model of Kahn[4] for inter-valence-band transitions. To get agreement with experimental results it was necessary to use two adjustable parameters plus a contribution from intraband free-carrier absorption. At 90 °K, WD identified the absorption as being due to neutral acceptors. No calculation of this absorption was made.

In this work, we present a model for absorption from neutral acceptors with applications to GaAs and GaP. The basic absorption mechanism is that of transitions of holes from the acceptor state to the sublevels of the valence band. In Sec. II we develop the necessary theory for absorption due to such transitions. Analytical calculations based on a simple band structure are performed and a comparison is made with the absorption from inter-valence-band transitions of free holes. In order to perform numerical calculations on GaAs and GaP we make approximations for the acceptor wave function. Here, we use a previously analyzed form for the ground-state wave function.[8] In Sec. III we give a description of the computation of the absorption in a model[2,7] which takes into account the detailed structure of the valence band and the optical matrix elements. In Sec. IV we present experimental results for the absorption of *p*-type GaAs together with the experimental results of WD for *p*-type GaP. The experimental data are compared with computed spectra for inter-valence-band absorption and for absorption from neutral acceptors.

II. ABSORPTION FROM NEUTRAL ACCEPTORS

A. Theory

At a sufficiently low temperature, the free holes in a semiconductor return to the acceptor ions and then neutralize these negatively charged states. The hole is in other words bound to the negative acceptor center. Let the valence-band Bloch functions with wave vector $\vec{k}$ be given by $\phi_{i,\vec{k}}(\vec{r}) = u_{i\vec{k}}(\vec{r})e^{i\vec{k}\cdot\vec{r}}$, where $i = 1, 2$, and 3 denote the heavy-hole band, light-hole band, and spin-orbit split-off band, respectively. $u_{i,\vec{k}}(\vec{r})$ are periodic func-