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RESONANT AMPLIFICATION AND COUPLING OF ACOUSTIC SURFACE WAVES WITH ELECTRONS DRIFTING ACROSS A MAGNETIC FIELD*

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We show that acoustic surface waves on a piezoelectric can be resonantly amplified by coupling to a weakly damped carrier surface wave on an adjacent semiconductor in which the conduction electrons are made to drift across an applied magnetic field. The conditions for such a wave-wave interaction are established, and calculations of the growth rate and its characteristics are given. Under the same conditions, reversing the direction of the magnetic field changes the amplifier to an almost loss-free, tunable, coupler.

The amplification of surface acoustic waves on a piezoelectric by drifting electrons in an adjacent semiconductor has been recently studied and demonstrated.[1] The mechanism of amplification in these studies, as in the well-known bulk amplifier,[2] is basically a nonresonant coupling of the acoustic wave to the electrons. That is, the electron system presents a small-signal, negative-real conductivity when its drift velocity v_0 exceeds the sound velocity v_a, but the wave is entirely established by the acoustic system. The consequence of this is that amplification occurs over a broad frequency range, with a frequency of maximum growth which is independent of v_0, and with growth rates which are proportional to the square of the effective electromechanical coupling constant, K^2. In contrast, as will be shown, in a resonant wave-wave interaction the frequency of maximum growth is directly controllable by v_0, and both the amplified bandwidth and the maximum growth rate are proportional to K. For such interactions to occur, the acoustic wave must be made to couple to a weakly damped wave carrying negative energy. For conduction electrons in a semiconductor,

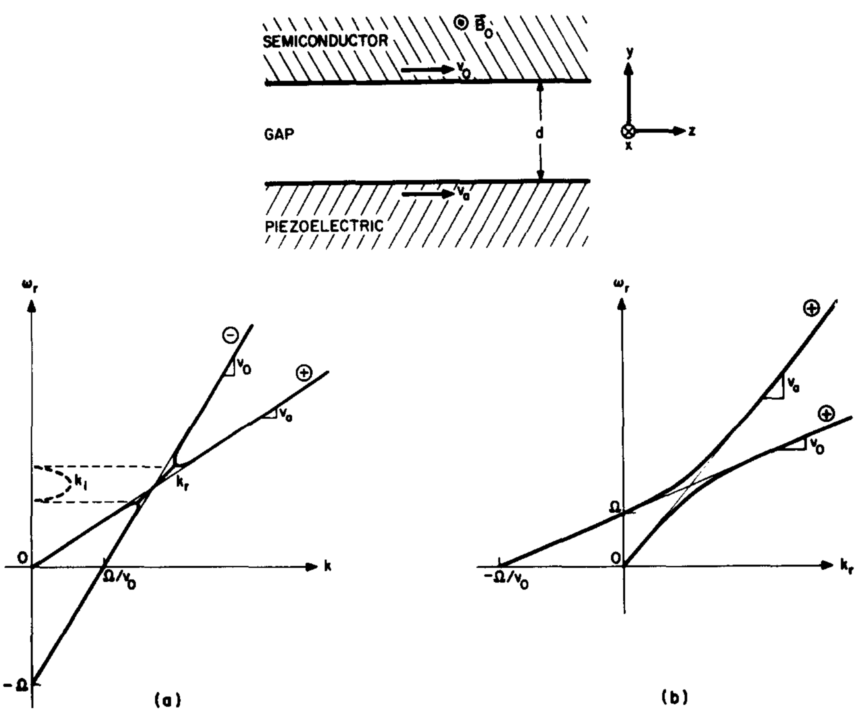

FIG. 1. Dispersion diagrams illustrating the resonant coupling of loss-free modes; (−) negative, and (+) positive energy modes; $\Omega = R\omega_c^* \, |b|$. (a) Magnetic field as shown ($B_0 < 0$), leading to amplification. (b) Magnetic field reversed ($B_0 > 0$), leading to coupling without amplification.

weakly damped surface waves exist when the electron drift velocity is across an applied magnetic field B_0 of magnitude such that $b \equiv \mu B_0 \gg 1$, where μ is the electron mobility.[3]

Consider the parallel-plane geometry of Fig. 1. We assume a weak piezoelectric on which a surface wave $\exp j(\omega t - kz)$ propagates in the z direction. We can then represent the effective permittivity for the piezoelectric by[4]

$$\epsilon_{pp} = \epsilon_p[1 + (\Delta v/v)(1 + \epsilon_0/\epsilon_p)k_{af}/(k - k_0)],$$

where ϵ_p is the permittivity in the absence of piezoelectricity,[5] $\Delta v/v$ is a parameter expressing the small piezoelectric coupling constant,[6] k_{af} is the wave number of the free acoustic surface wave in the absence of the semiconductor ($\epsilon_{pp} = -\epsilon_0$), $k_0 = k_{af}(1 + \Delta v/v)$ is the wave number when the electric field outside the piezoelectric is shorted out, and ϵ_0 is the permittivity in the gap. The semiconductor will be characterized by a single carrier (electron) dc conductivity σ_0, drift velocity v_0 across B_0, and permittivity ϵ_s. We use the small-signal, local, equations that describe the surface wave on the semiconductor in the absence of diffusion.[3] The effects due to diffusion will be described later. Assuming that the wave fields are electrostatic, $\bar{E} = -\nabla\Phi$, the potentials in the three regions of space are required to satisfy the continuity of E_z and D_y at $y = \pm d/2$, with proper account for the surface current at $y = d/2$. The detailed field analysis will be given elsewhere.[7] Here we discuss the resultant dispersion relation:

$$(k - k_a)(k - k_c) = -\frac{\Delta v}{v}\frac{G}{2}\frac{R\omega_c^*}{v_0}(b - j)k_{af}, \tag{1}$$

where

$$k_a = k_{af}\left[1 + \frac{\Delta v}{v}\frac{\epsilon_0(G_0 - 1)}{\epsilon_p + \epsilon_0 G_0}\right], \tag{2}$$

$$k_c = \epsilon/v_0 + (R\omega_c^*/v_0)(b - j), \tag{3}$$

$$R = \frac{\epsilon_s G_0(\epsilon_0 + \epsilon_p \tanh kd)}{(\epsilon_0 G_0 + \epsilon_p)(\epsilon_s + \epsilon_0 \tanh kd)}, \tag{4}$$

$$G = \frac{2\epsilon_0(\epsilon_p + \epsilon_0)(G_\infty - G_0)}{(\epsilon_p + \epsilon_0 G_0)(\epsilon_p + \epsilon_0 G_\infty)}, \tag{5}$$

$$G_0 = (\epsilon_s + \epsilon_0 \tanh kd)/(\epsilon_s \tanh kd + \epsilon_0), \tag{6}$$

$$G_\infty = \coth kd, \tag{7}$$

and $\epsilon_c^* = \sigma_0/\epsilon_s(1 + b^2)$. We note that for a given gap separation d, k_a is the propagation constant of the acoustic surface wave when $\sigma_0 = 0$, and k_c is the propagation constant of the carrier surface wave when $\Delta v/v = 0$. The dispersion relation, Eq. (1), represents the coupling of these waves, the right-hand side being the coupling coefficient which is a slowly varying function of k, and proportional to an effective electromechanical coupling constant $(\Delta v/v)G \equiv K^2$. Assuming that the acoustic surface wave is essentially undamped,[8] we find that Eq. (1) may exhibit two distinct types of interactions depending upon the damping rate of the carrier surface wave, which is given by k_{ci} — the imaginary part of Eq. (3).

For $|b| \gg 1$, we have, from Eq. (3), $k_{ci} \ll k_{ci}$ and the carrier surface wave is then only weakly damped. Under these conditions, resonant wave-wave interactions can occur at frequencies given by $k_a = \simeq k_{cr}$, as shown in Fig. 1. It can be shown[9] that for $b > 0$ the carrier surface wave has a negative small-signal energy, and hence will amplify the acoustic wave, as shown in Fig. 1(a). The frequency of maximum growth rate k_i occurs at $k_a = k_{cr}$ and is therefore tunable with v_0. The maximum growth rate and the amplified bandwidth are both proportional to the square root of the right-hand side of Eq. (1), and hence to $(G\Delta v/v)^{1/2} = K$. For $b < 0$, the carrier surface wave has a positive small-signal energy and the interaction with the acoustic wave is that of a coupler, as shown in Fig. 1(b). Signals can be coupled from the acoustic wave to the carrier wave, or vice versa; the coupler length for maximum signal transfer at $k_a = k_{cr}$ is also proportional to the square root of the right-hand side of Eq. (1).

For $|b| \sim 1$ or $|b| < 1$, Eq. (3) shows that the carrier surface wave is heavily damped ($k_{ci} \gtrsim k_{cr}$). Under these conditions, the interaction between the acoustic wave and carriers is nonresonant. Since the perturbation of the acoustic wave number is small in any case, we write approximately $(k - k_c) \approx (k_a - k_c) \approx (\omega/v_a - k_c)$ and obtain from Eq. (1)

$$\frac{k - k_a}{k_{af}} \simeq -\frac{\Delta v}{v}\frac{G}{2}\frac{(b - j)R\omega_c^*/\omega}{\eta - (b - j)R\omega_c^*/\omega}, \tag{8}$$

where $\eta = (v_0/v_a) - 1$. Equation (8) is precisely in the form of the well-known bulk-wave interaction,[2,10] where again $(\Delta v/v)G$ acts as an effective electromechanical coupling constant K^2, and $R\omega_c^*$ is the effective carrier relaxation frequency. The imaginary part of Eq. (8) gives the growth rate of the acoustic surface wave and is thus seen to be proportional to the real part of the effective small-signal conductivity of the electrons. For $b = 0$, we find the results previously obtained for $B_0 = 0$.[1] Equation (8) [but not Eq. (1)] has also been obtained recently by Greebe.[11] The growth rate k_i is now only proportional to $G(\Delta v/v) = K^2$, and the amplifier bandwidth is entirely determined by the

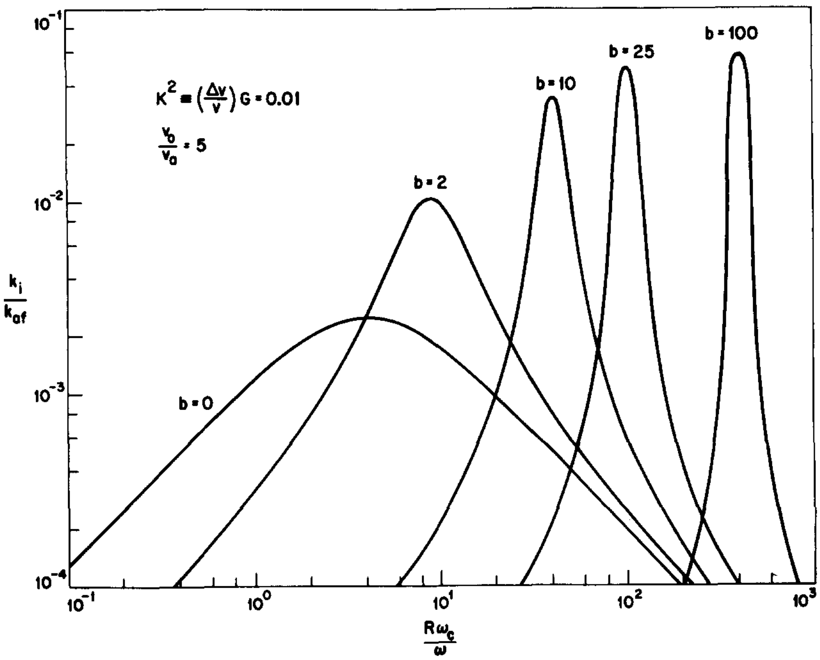

FIG. 2. Growth rate of acoustic surface wave k_i normalized to the free acoustic wave number k_{af}, as a function of the ratio of the effective carrier relaxation frequency $R\omega_c$ to the frequency ω.

frequency dependence of the carrier small-signal conductivity evaluated at the acoustic phase velocity.

For arbitrary values of b, both resonant and nonresonant interactions must be considered, and Eq. (1) must be solved without further approximations. Figures 2 and 3 show such solutions for the growth rate of the acoustic wave. The nature of the resonant wave-wave interactions is clearly evident for $b \gg 1$, where Eq. (8) is not applicable.

Finally, we consider the field analysis including the effect of diffusion on the carrier surface wave.[7]

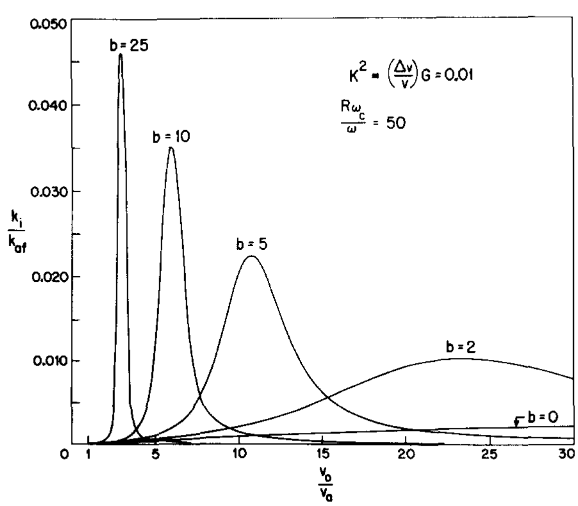

FIG. 3. Growth rate of acoustic surface wave k_i normalized to the free acoustic wave number k_{af}, as a function of the ratio of the electron drift velocity v_0 to the acoustic surface wave velocity v_a.

The small-signal electric potential inside the semiconductor must now be modified to include a part that derives from the bulk charge density, and the boundary condition of a surface current at $y = d/2$ is replaced by the vanishing of the normal component of the electron velocity.[12] For the resonant wave-wave interaction to dominate, it is sufficient to require that the damping of the carrier surface wave still be small. Including diffusion, we find that in addition to $|b| \gg 1$ we also require $|b| \times |j(\omega/\omega_D^*)/[\eta + j(\omega_c^*)]|^{1/2} \ll 1$, where $\omega_D^* = (1+b^2) \times v_a^2/\kappa T\mu/e$ is the diffusion frequency in the presence of a crossed magnetic field, and T is the electron temperature. Since the interaction frequency for $b \gg 1$ is $\omega \approx R\omega_c/b\eta$, the above condition for negligible damping due to diffusion becomes $\eta \gg (R\omega_c/b\omega_D)^{1/2}$, where $\omega_c = \sigma_0/\epsilon_s$ and $\omega_D = v_a^2/\kappa T\mu/e$.

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[1]Yu. V. Gulyaev and V. I. Pustovoil, Soviet Phys. JETP 20, 1508 (1965); K. A. Ingebritsen, ELAB Report TE 94, Norwegian Institute of Technology, Trondheim, Norway, 1967 (unpublished); J. H. Collins, K. M. Laken, C. F. Quate, and H. J. Shaw, Appl. Phys. Letters 13, 314 (1968).

[2]D. L. White, J. Appl. Phys. 33, 2547 (1962); A. R. Hutson, D. L. White, and J. H. McFee, Phys. Rev. Letters 7, 237 (1961).

[3]G. S. Kino, Appl. Phys. Letters 12, 312 (1968); R. Hirota, J. Phys. Soc. Japan 23, 798 (1968).

[4]K. A. Ingebritsen, J. Appl. Phys. 40, 2681 (1969); C. F. Quate, Microwave Lab. Report No. 1810, Stanford

University, 1969 (unpublished).

[5]In general, the piezoelectric is an anisotropic dielectric and $\epsilon_p = (\epsilon_{xx}\epsilon_{yy} - \epsilon^2_{xy})^{1/2}$.

[6]J. J. Campbell and W. R. Jones, IEEE Trans. Son. Ultrason. SU-15, 209 (1968); C. C. Tseng, J. Appl. Phys. 38, 4281 (1967). The parameter $\Delta v/v$ is the fractional change in the surface wave velocity when a shorting plane is placed close to, but not in mechanical contact with, the piezoelectric.

[7]A. Bers and B. E. Burke (unpublished).

[8]To account for a weakly damped acoustic wave, we need only add a small imaginary part to k_a.

[9]The time-averaged energy per unit length in z is proportional to $\partial D(\omega,k)/\partial\omega$, where $D(\omega,k)$ is the dispersion relation for the loss-free carrier surface wave in the absence of the piezoelectric.

[10]M. C. Steele, RCA Rev. 28, 58 (1967); G. S. Kino and R. Route, Appl. Phys. Letters 11, 312 (1967); A. Bers, MIT Res. Lab. Electron., Quarterly Prog. Rept. 88, 204 (1968).

[11]C. A. A. J. Greebe (unpublished).

[12]B. E. Burke, Microwave Lab. Report 1738, Stanford University, 1967 (unpublished).

LOW-POWER QUASI-cw RAMAN OSCILLATOR

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Quasi-cw stimulated Raman emission in the visible has been obtained from an oscillator cavity with a pump-power input of less than 5 W. This large reduction of oscillator threshold is achieved with the use of a liquid-core optical waveguide structure. Significant conversion of pump to Stokes light has been observed. Extension of the system to continuous operation and to the study of other nonlinear effects is suggested.

This letter reports the first observation of stimulated Raman emission from an oscillator cavity under conditions of low-power quasicontinuous excitation. In contrast to previous Raman lasers which have required pump powers in excess of 10^4 W,[1] the oscillator described here operates with an input power of less than 5 W. The key to this low-power operation is confinement of both the pump and Stokes beams in a liquid-core optical waveguide. It is the maintenance of high optical power densities over relatively long lengths of guide that dramatically increases the gain per pass in this type of oscillator. In our experiments the waveguides are fabricated by filling small-bore glass tubing with carbon disulfide (CS_2). The guide is pumped longitudinally with a repetitively pulsed argon-ion laser operating at $\lambda_p = 5145$ Å. Raman emission occurs at a Stokes wavelength $\lambda_s = 5325$ Å. This observed Stokes frequency shift is characteristic of the 656-cm^{-1} line in CS_2.

Operation of the oscillator is illustrated schematically in Fig. 1. A magnified view of the guide

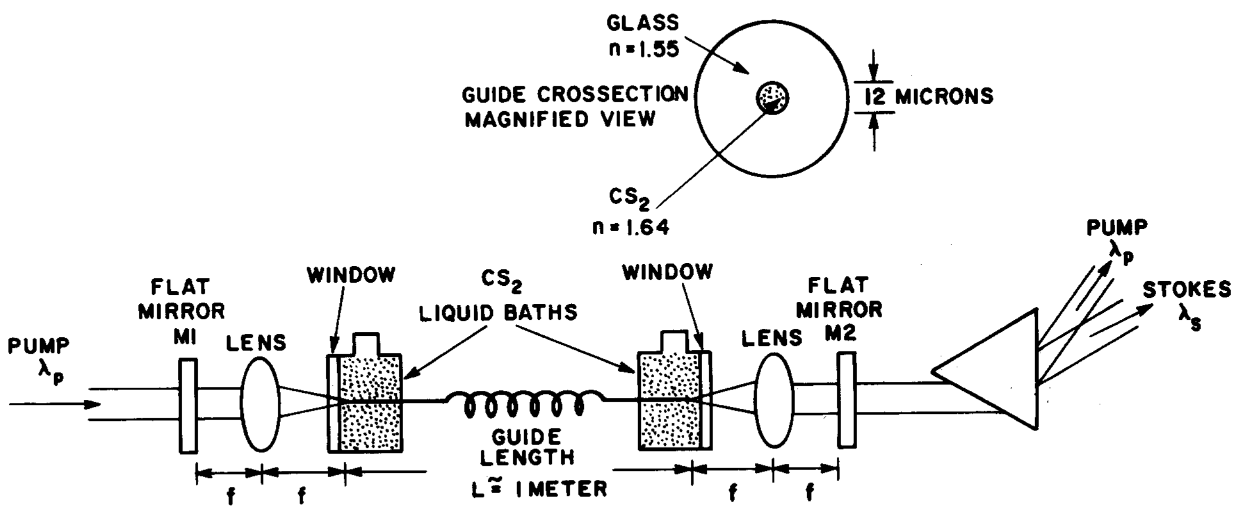

FIG. 1. Experimental arrangement for the observation of Raman oscillation in a liquid-core optical waveguide resonator.