PHYSICAL REVIEW B VOLUME 18, NUMBER 10 15 NOVEMBER, 1978

Magnetoplasma polaritons at the interface between a semiconductor and a metallic screen

P. Halevi and C. Guerra-Vela

Instituto de Ciencias de la Universidad Autónoma de Puebla, Apartado Postal J-48, Puebla, Puebla, México

(Received 27 June 1978)

An interface between a semiconductor and a highly conducting metallic screen supports polaritons that propagate at a right angle to a static magnetic field lying in the plane of the interface (Voigt geometry). In case of negligible plasmon-phonon interaction, there is a propagation window below the cyclotron frequency and another one above the hybrid cyclotron-plasma frequency. For a given orientation of the magnetic field, these modes propagate in opposite directions. In case of a *polar* semiconductor, there are three propagation windows whose limiting frequencies depend on the cyclotron, plasma, transverse-phonon, and longitudinal-phonon frequencies. These modes also exhibit nonreciprocality with respect to their direction of propagation. Effects of damping and of the finite conductivity of the metallic screen, and the possibility of experimental detection of the predicted modes are discussed.

I. INTRODUCTION

Surface polaritons in semiconductors have been extensively studied in the course of the past decade. In particular, the effect of a static magnetic field was investigated in detail, and the reader is referred to review articles[1-4] on the subject. The magnetic field is usually taken to be either perpendicular or parallel to the surface. In the latter case the most frequently studied configurations are propagation along the magnetic field (Faraday geometry) and at a right angle to the field (Voigt geometry). A detailed study of the three geometries was made by Wallis *et al.*[5] In some publications[1,2,4-6] it was assumed that the effect of phonons is unimportant; this is justified provided that the plasma and cyclotron frequencies of the charge carriers and the frequency of the wave are much higher than the longitudinal-phonon frequency. Other works[3,7,8] allow for the interaction between plasmons and phonons, resulting in an increased number of polariton modes.

Whether the effect of the phonons is included or not included in the description of the polar semiconductor, the presence of an applied magnetic field causes various qualitative changes. (i) The modes that exist in the absence of a magnetic field are shifted in frequency, the shift being of the order of the cyclotron frequency. (ii) An originally isotropic semiconductor becomes highly anisotropic, i.e., the properties of the modes and their very existence depend on the relative directions of the magnetic field, the direction of propagation, and the surface normal. (iii) The polariton modes depend not only on the direction of the magnetic field, but also on its sense. In other words, upon reversing the field the original modes disappear and the new ones appear. This behavior is often referred to as "nonreciprocal." (iv) In the Voigt geometry certain modes stop at a finite value of the wave vector. Therefore they are not true excitations of the plasma at the crystalline surface. They are called "virtual excitations" or "photon-driven" surface plasmons. (v) If the magnetic field is perpendicular to the surface and also in the Faraday geometry the surface mode is composed of a linear combination of two inhomogeneous waves which, in general, have different decay constants. This feature gives rise to a variety of unusual surface modes ("generalized,"[5,9] "leaky,"[5,9] and "degenerate"[10,11]), in addition to the standard ones. (vi) With the exception of the Voigt geometry, the plane of polarization of the electric field does not coincide with the sagittal plane. As a consequence, a combination of s- and p-polarized light is necessary in order to excite the surface-polariton modes by optical techniques.

Such a technique-attenuated total reflection (ATR) was employed by Palik *et al.*[9,12] in order to excite surface polaritons in n-type InSb in the Faraday geometry. Experiments were also performed in the much simpler case of the Voigt geometry by Hartstein *et al.*[13-15] These authors confirmed the existence of nonreciprocal behavior and virtual excitations for InSb with a high concentration of electrons (negligible interaction with phonons)[13] and with a low concentration of electrons (plasmon-phonon interaction).[14]

The effect of an electron-density profile at the semiconducting surface was investigated by Wallis *et al.*[16] In the presence of such a profile nonlocal effects[17] play an important role. This is because standing waves of small wavelength may be excited in the "step" of the density profile.

The problem of polariton modes propagating at the interface between two magnetoplasmas does not lend itself to an easy solution. The case of the Voigt geometry was recently studied in detail by

Uberoi and Rao.[18,20] They derived a general, however, implicit dispersion relation [see Eq. (1) of Ref. 19], which was subsequently applied to the special cases that both media are (a) plasmas[18] (with the effect of ions included); (b) semiconductors[19] (neglecting the effect of phonons); and (c) polar semiconductors.[20] The authors analyzed the problem in terms of a set of parameters relevant to the interface. Unfortunately the analysis is quite complicated; for instance, in the case of an interface between two plasmas the roots of the dispersion relation are discussed in terms of a diagram divided into no less than 25 regions.[18] For an interface between two semiconductors there are—depending on the orientation of the magnetic field—one or two authentic surface modes.[19] In addition to these there may appear in the spectrum gaps associated with virtual excitations. An interface between two *polar* semiconductors may support as many as seven modes.[20]

II. POLARITONS AT THE INTERFACE BETWEEN A POLAR SEMICONDUCTOR AND A METALLIC SCREEN IN THE VOIGT GEOMETRY

In this paper we discuss a particularly simple case of an interface between two conductors, whose mathematical analysis and physical interpretation are straightforward and transparent. One conductor is a polar semiconductor and the other is a highly conducting metallic screen. The wave fields cannot penetrate into the latter medium; in particular, the tangential component of the electric field and the normal component of the magnetic field vanish. Then, as a result of the continuity of these components across the boundary, the electric field $\vec{E}$ in the semiconducting medium must be perpendicular to the interface and the magnetic field $\vec{B}$ must be parallel to the interface. Subject to these boundary conditions we solve Maxwell's equations for the polar semiconductor in the presence of a static magnetic field

$$\vec{\nabla}\times\vec{E} = -\frac{1}{c}\frac{\partial\vec{B}}{\partial t}, \tag{1}$$

$$\vec{\nabla}\times\vec{B} = \frac{1}{c}\frac{\partial}{\partial t}(\underline{\epsilon}\cdot\vec{E}). \tag{2}$$

Here $\underline{\epsilon}$ is the magnetoplasma tensor of the medium. Eliminating $\vec{B}$ we easily arrive at the following wave equation[21]

$$\nabla^2\vec{E} - \vec{\nabla}(\vec{\nabla}\cdot\vec{E}) - \frac{1}{c^2}\frac{\partial^2}{\partial t^2}\underline{\epsilon}\cdot\vec{E} = 0. \tag{3}$$

For $\vec{E}\propto\exp[i(\vec{q}\cdot\vec{r}-\omega t)]$ the last equation becomes

$$q^2\vec{E} - \vec{q}(\vec{q}\cdot\vec{E}) - q_0^2\underline{\epsilon}\cdot\vec{E} = 0 \tag{4}$$

where $q_0 = \omega/c$ is the vacuum wave vector. This equation is valid for an arbitrary direction of the static magnetic field $\vec{B}_0$. If we choose the z axis along $\vec{B}_0$ the dielectric tensor is given by[22]

$$\underline{\epsilon} = \begin{bmatrix} \epsilon_{xx} & -\epsilon_{yx} & 0 \\ \epsilon_{yx} & \epsilon_{xx} & 0 \\ 0 & 0 & \epsilon_{zz} \end{bmatrix}, \tag{5}$$

where

$$\epsilon_{xx} = \epsilon_\infty\left(\frac{\omega^2-\omega_L^2}{\omega^2-\omega_T^2} - \frac{\omega_p^2}{\omega^2-\omega_c^2}\right), \tag{6}$$

$$\epsilon_{yx} = -i\,\epsilon_\infty\frac{\omega_c\,\omega_p^2}{\omega(\omega^2-\omega_c^2)}, \tag{7}$$

$$\epsilon_{zz} = \epsilon_\infty\left(\frac{\omega^2-\omega_L^2}{\omega^2-\omega_T^2} - \frac{\omega_p^2}{\omega^2}\right). \tag{8}$$

Here ω_T and ω_L are the transverse- and longitudinal-phonon frequencies, respectively, $\omega_p = (4\pi ne^2/m^*\epsilon_\infty)^{1/2}$ is the plasma frequency of the charge carriers, $\omega_c = eB_0/m^*c$ is the cyclotron frequency, and ϵ_∞ is the high-frequency dielectric constant. We have neglected the damping frequencies of both lattice vibrations and plasma waves.

We shall now limit the static magnetic field to a direction parallel to the interface. If we choose the y axis normal to the interface then our boundary conditions imply that the only nonzero component of the electric field is E_y, i.e., $E_x = E_z = 0$. All the nonvanishing terms in Eq. (4) are proportional to E_y, which may be cancelled out. Then the x, y, and z components of Eq. (4) reduce to the following equations:

$$q_x q_y = q_0^2\epsilon_{yx}, \tag{9}$$

$$q_x^2 + q_z^2 = q_0^2\epsilon_{xx}, \tag{10}$$

$$q_y q_z = 0. \tag{11}$$

By Eq. (9) $q_y \neq 0$ and therefore by Eq. (11), $q_z = 0$. This means that there is no wave propagation along the magnetic field. In other words, the Faraday configuration does not support a *single*-plane-wave solution. Equations (9) and (10) become[23]

$$q_x = q_0\sqrt{\epsilon_{xx}}, \tag{12}$$

$$q_y = q_0\epsilon_{yx}/\sqrt{\epsilon_{xx}}. \tag{13}$$

The conditions for propagation of an interface polariton are that q_x is real (positive or negative) and that iq_y is real and negative. The latter condition ensures that the wave decays exponentially away from the interface in the semiconductor ($y>0$). By Eq. (6) ϵ_{xx} is real: therefore, by Eq. (12), q_x is real whenever $\epsilon_{xx}>0$. The wave may propagate in the positive or negative direction of the x axis ($q_x>0$ or $q_x<0$) depending on the sign of the square root in Eq. (12). If $q_x>0$ then, by Eq. (9), iq_y is real and negative for $\omega<\omega_c$. On the other hand, if $q_x<0$, then iq_y is real and negative for $\omega>\omega_c$. We may then conclude that *the interface supports polaritons whenever* $\epsilon_{xx}>0$. *They propagate in the positive (negative) x directions for* $\omega<\omega_c$ ($\omega>\omega_c$). This behavior corresponds to the nonreciprocal nature of interface polaritons in the presence of a static magnetic field.

A very simple physical picture of the wave emerges (see Fig. 1). It propagates at right angles to the applied magnetic field (Voigt geometry) with a phase velocity $\omega/q_x=c/\sqrt{\epsilon_{xx}}$ [see Eq. (12)].[24] It decays exponentially into the semiconducting medium with a decay constant given by $|q_y|$ [see Eq. (13)].[25] The wave is linearly polarized with its electric field perpendicular to the interface. It follows from Eq. (1) that the magnetic field of the wave is parallel to $\vec{B}_0$. Using Eqs. (1) and (12) we get $B_z=\sqrt{\epsilon_{xx}}\,E_y$.

ϵ_{xx} [see Eq. (6)] has two poles, at ω_T and at ω_C. It also has two zeros, ω_1 and ω_2, which are

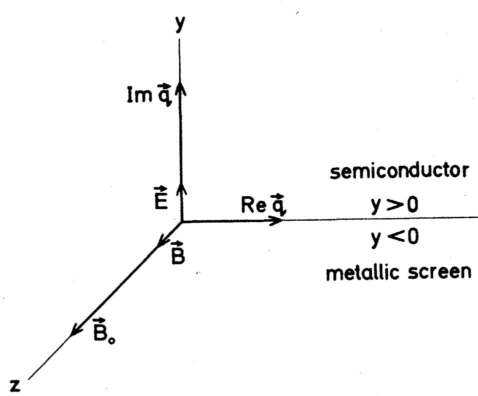

FIG. 1. Voigt geometry: the interface wave propagates in the direction $\mathrm{Re}\,\vec{q}$, at a right angle to the static magnetic field $\vec{B}_0$; both vectors are parallel to the interface. The wave is linearly polarized, with its electric field $\vec{E}$ perpendicular to the interface and its magnetic field $\vec{B}$ parallel to the static field $\vec{B}_0$. The decay into the semiconductor is given by $\mathrm{Im}\,\vec{q}$, while the decay into the metallic screen is negligible. Note that the electric field acting on an electron causes a drift velicity $\vec{v}_d$ in the $-y$ direction. Then there is an additional force, $-e\vec{v}_d\times\vec{B}_0$, in the $+x$ direction. This preferential direction is responsible for the nonreciprocal nature of propagation.

the solutions of the equation

$$\omega^4-(\omega_L^2+\omega_P^2+\omega_C^2)\omega^2+(\omega_T^2\omega_P^2+\omega_L^2\omega_C^2)=0\,. \tag{14}$$

Now $\omega_T<\omega_L$, and we may choose $\omega_1<\omega_2$. Taking in account that $\epsilon_{xx}(0)$ and $\epsilon_{xx}(\infty)$ are positive quantities it is a simple matter to establish that

$$\min(\omega_T,\ \omega_c)<\omega_1<\max(\omega_T,\ \omega_c)<\omega_2 \tag{15}$$

holds for an arbitrary set of characteristic frequencies ω_T, ω_L, ω_p, and ω_c. Here $\min(\omega_T,\omega_c)$ is the smaller one among the frequencies ω_T and ω_c and $\max(\omega_T,\omega_c)$ is the greater one among these frequencies. In Fig. 2(a) we present a schematic plot of the function $\epsilon_{xx}(\omega)$. It is positive in the regions indicated by heavy lines on the frequency axis. Thus there are three "propagation windows" for interface polaritons, defined as follows:

$$\text{I}\quad \omega<\min(\omega_T,\ \omega_c)\,; \tag{16}$$

$$\text{II}\quad \omega_1<\omega<\max(\omega_T,\ \omega_c)\,; \tag{17}$$

$$\text{III}\quad \omega>\omega_2\,. \tag{18}$$

In these regions the dispersion relation of the modes is given by Eq. (12) and is schematically drawn in Fig. 2(b). It is remarkable that this general behavior is valid for an arbitrary amount of doping and arbitrary $\vec{B}_0$.

Now in region I ω is always smaller than ω_c. Therefore, the low-frequency mode in Fig. 2(b) may propagate only in the positive x direction ($q_x>0$). On the other hand, in region III ω is always greater than ω_c. Thus the high-frequency mode in Fig. 2(b) may propagate only in the negative x direction ($q_x<0$). In region II $\omega<\omega_c$ when $\omega_T<\omega_c$, and vice versa. This means that mode II propagates in the positive (negative) x direction when $\omega_T<\omega_c$ ($\omega_T>\omega_c$). This nonreciprocal behavior is a characteristic feature of magnetoplasma surface and interface polaritons. It has been both predicted[7,26] and observed[13,14] for semiconducting surfaces in the Voigt geometry. It has been also predicted for interfaces between two semiconductors in the same geometry.[19,20] The effect may be observed experimentally by reversing the direction of the applied magnetic field, thereupon observing polaritons which were formerly absent.

The modes of Fig. 2(b) should be readily observable by attenuated total reflection (ATR) techniques. We stress that in the present case there should be no "air gap": a semiconducting film of optimum thickness should be "sandwiched" between a metallic screen and a high-index prism.[27] Otherwise polariton modes at the free semiconducting surface would be excited in addition to the modes at the semiconductor-screen interface. These modes would be coupled and would lie to the right of the vacuum light line, unlike the polariton modes in Fig. 2(b). Such coupled modes in the absence of an applied field were predicted[28] and observed by López-Ríos *et al.*[29]

Recently, Halevi and Hernández-Cocoletzi[30] calculated the reflectivity spectrum of a prism-metal-metal configuration. Well-defined minima, which derive from the excitation of plasmons at the bimetallic interface, were found. These interface plasmon-polariton modes were predicted[31] and recently reviewed[32] by the first author. We expect that the strong excitation will persist in the presence of an applied magnetic field which, of course, must be perpendicular to the sagittal plane.

III. DAMPING EFFECTS

In Sec. II we neglected the effects of damping. If we allow for finite, phenomenological damping

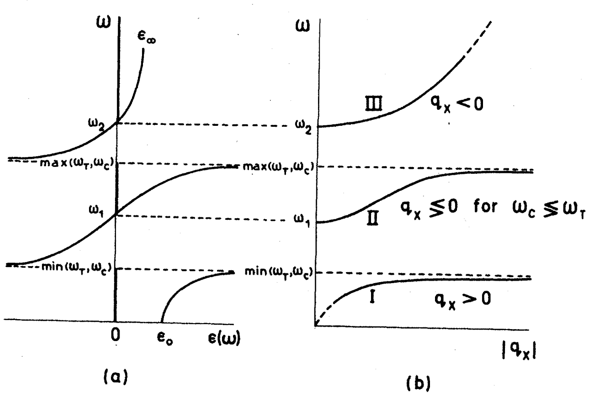

FIG. 2. (a) Element ϵ_{xx} of the dielectric tensor vs ω, for a polar semiconductor, Eq. (6). This schematic representation is valid for an arbitrary configuration of the characteristic frequencies ω_T, ω_L, ω_c, and ω_p. Whenever $\epsilon_{xx} > 0$ a propagation window for interface polaritons exists, as marked by heavy lines along the frequency axis. (b) The polariton modes of an interface composed of a polar semiconductor and a metallic screen, corresponding to the propagation windows of (a). Mode I propagates in the $+x$ direction, mode III propagates in the $-x$ direction, and mode II propagates in the $+x$ $(-x)$ direction for $\omega_c > \omega_T$ $(\omega_c < \omega_T)$. The reason for this nonreciprocal behavior is explained in the legend in Fig. 1. The dashed lines indicate that our model breaks down at very low frequencies $(\omega \sim \nu)$ and at very high frequencies $(\omega \sim \omega_p')$. Mode III actually terminates at the resonant frequency given by Eq. (22).

frequencies for the phonon and plasmon terms in Eqs. (6)–(8) then, in particular, ϵ_{xx} becomes complex. Assuming that its imaginary part ϵ''_{xx} is much smaller than its real part ϵ'_{xx} Eq. (6) becomes

$$\begin{aligned} q_x &= q_0(\epsilon'_{xx} + i\,\epsilon''_{xx})^{1/2} \\ &\cong q_0\sqrt{\epsilon'_{xx}} + i\,q_0\epsilon''_{xx}/2\sqrt{\epsilon'_{xx}}\,. \end{aligned} \tag{19}$$

For spatial damping ω is real and $q_x \equiv q'_x + i\,q''_x$ is a complex quantity. The dispersion relation, ω vs q'_x, will suffer certain changes. The most important is that the propagation window I [Eq. (16)] will not extend to the very-low-frequency region as in Fig. 2. In fact, propagation will be limited to frequencies $\omega \gtrsim \nu$, where ν is the plasmon damping frequency (it is typically much larger than the phonon damping frequency). We also expect strong damping effects in the vicinity of the frequencies ω_1, ω_2, ω_T, and ω_c. As a result, the propagation windows II and III will not be sharply delineated any more. The wave vector q'_y will not go to zero close to ω_1 and ω_2. Neither will it become infinite as $\omega \to \omega_T$ and $\omega \to \omega_c$. As usual with the case of spatial damping, the dispersion relation should bend back as these frequencies are approached.

The second term in Eq. (19) describes the attenuation of the wave along its direction of propagation. The propagation distance is given by

$$L = 1/2q''_x = \sqrt{\epsilon'_{xx}}/q_0\epsilon''_{xx}\,. \tag{20}$$

IV. CASE OF NEGLIGIBLE PLASMON-PHONON INTERACTION

It is interesting to see what happens in a case of experimental interest, namely, that $\omega_L \ll \min(\omega_c,\ \omega_p)$. Then the interaction between plasmons and phonons may be neglected, and we may take $\omega_T = \omega_L = 0$ in the frequency region $\omega \gg \omega_L$. The solutions of Eq. (14) are $\omega_1 = 0$ and $\omega_2 = (\omega_p^2 + \omega_c^2)^{1/2}$. Then the propagation windows I and II collapse into a single window, with no stopband between $\min(\omega_T,\ \omega_c)$ and ω_1. The schematic dispersion relation is drawn in Fig. 3. In this case the lower- (higher-) frequency mode always propagates in the positive (negative) direction of the x axis.

In the absence of a static magnetic field $\omega_c = 0$, and the lower-frequency mode in Fig. 3 disappears. The higher-frequency mode now starts at ω_p. This is the polariton mode at the interface between two conductors, predicted by the first author.[31] However this model has a limiting frequency, given by

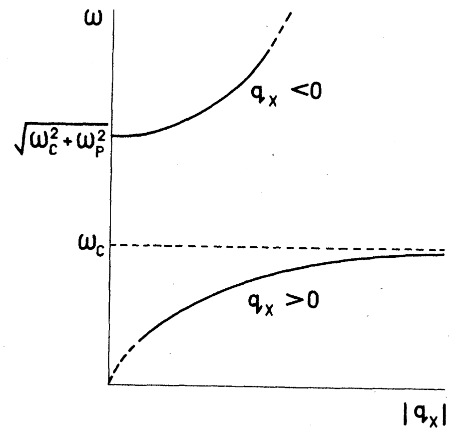

FIG. 3. Polariton modes of an interface composed of a semiconductor and a metallic screen in the case of negligible plasmon-phonon interaction. [ϵ_{xx} is given by Eq. (6) with $\omega_T=\omega_L=0$.] The meaning of the dashed lines is explained in the legend to Fig. 2(b).

$$\omega_s = \left(\frac{\epsilon_\infty\omega_p^2+\epsilon'_\infty\omega_p'^2}{\epsilon_\infty+\epsilon'_\infty}\right)^{1/2} \tag{21}$$

The plasma frequency ω'_p of a "highly conducting metallic screen" is much larger than ω_p of the semiconductor, and $\epsilon'_\infty=1$. Then Eq. (21) may be approximated by

$$\omega_s \cong \omega'_p/(\epsilon_\infty+1)^{1/2}. \tag{22}$$

Now our model implicitly assumes that $\omega'_p\to\infty$. Therefore it is not surprising that $\omega\to\infty$ as $q_x\to\infty$ for the high-frequency modes in Figs. 2(b) and 3. It becomes clear that this behavior is a manifestation of the limitations of our model. For a metallic medium with a *finite* plasma frequency ω'_p, the high-frequency modes in Figs. 2(b) and 3 should terminate at some cutoff frequency. This frequency depends, in principle, on the parameters which characterize the two conductors in contact; however, it is approximately given by Eq. (22) provided that ω'_p is much larger than $\max(\omega_L,\ \omega_p,\ \omega_c)$.[33]

Another effect of a finite ω'_p is that the polariton falls off exponentially into the metallic screen. Its decay distance is of the order of c/ω'_p, rather than zero.

We note that the modes discussed in this work are of the standard interface polariton type. Thus we find no virtual (photon-driven) modes. This may be readily understood from a simple criterion for the existence of these modes, given by Uberoi and Rao.[19] The criterion is that the inequalities $\omega_p<\omega'_p$ and $\epsilon_\infty<\epsilon'_\infty$ must be simultaneously satisfied. In the present case, however, the second inequality does not hold ($\epsilon_\infty>1$ and $\epsilon'_\infty=1$).

V. CASE OF B_0 PERPENDICULAR TO THE SEMICONDUCTING SURFACE

Up to this point the applied magnetic field was constrained to lie in the plane of the interface. Now we discuss the case that it is perpendicular to the interface. Therefore $\vec{B}_0 \parallel \vec{E} \parallel \hat{z}$, where $\hat{z}$ is a unit vector normal to the interface. Equation (4) must be solved subject to the boundary conditions $E_x=E_y=0$. It is straightforward to show that

$$q=q_0\sqrt{\epsilon_{zz}},\quad q_z=0. \tag{23}$$

Thus the wave propagates in an arbitrary direction in the plane of the interface and does not decay at all into the semiconducting medium. This is simply a bulk polariton, with phonon-plasmon interaction. There are two propagation windows: one between ω' and ω_T, and the other above ω'', where ω' and ω'' are the zeros of $\epsilon_{zz}(\omega)$.

The properties of the wave are independent of the magnetic field. This is explained by the fact that the charge carriers are accelerated in the direction of $\vec{B}_0$ (because $\vec{E}\parallel\vec{B}_0$) and, therefore, experience no Lorentz force.

ACKNOWLEDGMENT

One of the authors (C. G. V.) acknowledges support from Consejo Nacional de Ciencia y Tecnologia.

[1]K. W. Chiu and J. J. Quinn, Nuovo Cimento B 10, 1 (1972).
[2]J. J. Quinn and K. W. Chiu, in *Proceedings of the Taormina Research Conference on the Structure of Matter-Polaritons*, 1972 (Pergamon, Oxford, 1974), p. 259.
[3]K. L. Kliewer and R. Fuchs, Adv. Chem. Phys. 27, 355 (1974).
[4]A. Otto, in *Optical Properties of Solids, New Developments*, edited by B. O. Seraphin (North-Holland, Amsterdam, 1975).
[5]R. F. Wallis, J. J. Brion, E. Burstein, and A. Hartstein, Phys. Rev. B 9, 3424 (1974).
[6]A. B. Davydov and V. A. Zakharov, Fiz. Tverd. Tela 17, 201 (1975) [Sov. Phys. Solid State 17, 117 (1975)].
[7]K. W. Chiu and J. J. Quinn, Phys. Rev. Lett. 29, 600 (1972).
[8]J. J. Brion, R. F. Wallis, A. Hartstein, and E. Burstein, Surf. Sci. 34, 73 (1973).
[9]E. D. Palik, R. Kaplan, R. W. Gammon, H. Kaplan,

R. F. Wallis, and J. J. Quinn, Phys. Rev. B 13, 2497 (1976).
[10]P. G. Flahive and J. J. Quinn, Phys. Rev. Lett. 9, 586 (1973).
[11]U. J. Rao and C. Uberoi, Phys. Rev. B 18, 2941 (1978).
[12]E. D. Palik, R. Kaplan, R. W. Gammon, H. Kaplan, J. J. Quinn, and R. F. Wallis, Phys. Lett. A 45, 143 (1973).
[13]A. Hartstein and E. Burstein, Solid State Commun. 14, 1223 (1974).
[14]A. Hartstein, E. Burstein, E. D. Palik, R. Kaplan, R. W. Gammon, and B. W. Henvis, in *Proceedings of the International Conference on the Physics of Semiconductors*, Stuttgart, 1974 (Teubner, Stuttgart, 1974), p. 546.
[15]A. Hartstein, E. Burstein, E. D. Palik, R. W. Gammon, and B. W. Henvis, Phys. Rev. B 12, 3186 (1975).
[16]R. F. Wallis, J. J. Brion, E. Burstein, and A. Hartstein, in *Proceedings of the Eleventh International Conference on the Physis of Semiconductors* (PWN-Polish Scientific, Warsaw, 1972), p. 1448.
[17]A. Eguiluz and J. J. Quinn, Phys. Rev. B 10, 4299 (1976); A. D. Boardman and E. D. Isaac, Solid State Commun. 18, 97 (1976); G. C. Aers and A. D. Boardman, J. Phys. C 11, 945 (1978).
[18]C. Uberoi and U. J. Rao, Plasma Phys. 17, 659 (1975).
[19]C. Uberoi and U. J. Rao, Surf. Sci. 66, 210 (1977).
[20]U. J. Rao and C. Uberoi, Phys. Lett. A (to be published).
[21]It follows from Eq. (2) that $\vec{\nabla}\cdot\vec{D}=\vec{\nabla}\cdot(\vec{\epsilon}\cdot\vec{E})=0$. However, $\vec{\nabla}\cdot\vec{E}\neq 0$, which is a consequence of the anisotropy caused by the static magnetic field. In fact, it follows from Eq. (13) and the considerations preceding it that $\vec{\nabla}\cdot\vec{E}=iq_yE_y=iE_yq_0\epsilon_{yx}/\sqrt{\epsilon_{xx}}$.
[22]See, for example, Ref. 9. We define ω_c as a positive quantity for an n-type semiconductor ($e>0$).
[23]Equation (12) may be also derived from the general dispersion relation of Uberoi and Rao (Ref. 19) by taking the limit that the plasma frequency of the metallic medium goes to infinity.
[24]It is interesting that the phase velocity does not depend on the off-diagonal element ϵ_{yx}. By way of comparison, the phase velocity of ***bulk*** polaritons in the Voigt geometry is $c/\sqrt{\epsilon_V}$, where $\epsilon_V=\epsilon_{xx}+\epsilon_{yx}^2/\epsilon_{xx}$. In fact, from Eqs. (12) and (13) we find $q_x^2+q_y^2=q_0^2\epsilon_V$.
[25]We assumed that the wave fields do not penetrate into the metallic screen, corresponding to a vanishing decay constant. If the metallic screen has a finite plasma frequency ω_p', the decay constant is approximately c/ω_p'.
[26]J. J. Brion, R. F. Wallis, A. Hartstein, and E. Burstein, Phys. Rev. Lett. 28, 1455 (1972).
[27]This corresponds to the "Otto Configuration" because the semiconductor is surface-wave inactive, while the metallic screen is surface-wave active. Another possibility is the "Kretschmann configuration" with a metallic screen of thickness of the order of (Ref. 25) c/ω_p' sandwiched between the prism and a semiconducting substrate. This configuration actually has the advantage that the optimum thickness of the film is practically independent of the frequency.
[28]T. López-Rios, Opt. Commun. 17, 342 (1976).
[29]T. López-Rios, F. Abeles, and G. Vuye, J. Phys. (Paris) (to be published).
[30]P. Halevi and G. Hernández-Cocoletzi, Phys. Rev. B 18, 590 (1978).
[31]P. Halevi, Phys. Rev. B 12, 4032 (1975).
[32]P. Halevi, Surf. Sci. 76, 64 (1978).
[33]A cubic equation for the resonant frequencies, derived by Uberoi and Rao (Ref. 19), reduces to Eq. (22) for very high values of ω_p'.