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Stimulated Raman scattering from plasma modes in magnetoactive semiconductors

S GUHA and N APTE*
Department of Physics, Ravishankar University, Raipur 492 010, India
*Present Address: School of Studies in Physics, Vikram University, Ujjain 456 331, India

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Abstract. Stimulated scattering off electron plasma mode is investigated analytically for the case when the pump wave is an intense circularly polarised electromagnetic wave propagating parallel to a homogeneous dc magnetic field in an isotropic semiconductor-plasma. The threshold electric field of the pump necessary for the stimulated Raman scattering and the growth rate of the parametrically unstable mode have been obtained for two cases (i) $B_0=0$ and (ii) $B_0\neq 0$. It is seen that the magnetic field does not significantly affect the threshold electric field as well as the growth rate provided the cyclotron frequency is small compared to the frequency of the pump wave. The threshold conditions are also found to be insensitive to the electron thermal velocity.

Keywords. Stimulated Raman scattering; parametric instability; piezoelectric semiconductor; plasma modes.

1. Introduction

There has been much interest in the scattering of intense electromagnetic radiation in a plasma resulting from the decay of the incident wave into another electromagnetic wave and either a Langmuir wave (stimulated Raman scattering, SRS) or an acoustic wave (stimulated Brillouin scattering). SRS provides useful information about the spectrum and nature of elementary excitations in a plasma. Originally the technique was applied to gaseous plasma (Kunze *et al* 1964). Several theoretical studies of stimulated Raman scattering (SRS) in an infinite homogeneous isotropic gaseous plasma have been reported recently (*e.g.* Drake *et al* 1974; Lashmore-Davies 1975; Fuchs 1976; Sodha *et al* 1976). The stimulated Raman backscattering of a circularly polarized electromagnetic wave propagating along an uniform static magnetic field has been investigated by Maraghechi and Willett (1979). The phenomenon of SRS has been used to study plasmas in semiconductors as well (Wolff 1970; Patel and Shaw 1970). Many theoretical treatments considered either the photon-system coupling with a single component one-band electron gas, omitting all phonon effects (Wolff 1970) or considered undoped semiconductors in which no conduction electron is present. Foo and Tzoar (1970) calculated the Raman scattering cross-section in the forward direction in solid-state plasma in a magnetic field. However, the threshold electric field required for the onset of SRS from electron plasma mode and the growth rate of the unstable mode well above the threshold as well as their dependence on the various parameters of the system have not been studied in solid-state plasma.

In this paper we report our investigations on the SRS of a circularly polarized electromagnetic wave due to the excitation of electron plasma waves propagating along a uniform static magnetic field in a semiconductor-plasma. The electron concentration in the semiconductor is chosen suitably so as to neglect the effect of optical phonons on SRS. The stimulated Raman forward as well as backscattering of a forward travelling pump wave have been studied incorporating its (pump) spatial dependence. The coupled mode theory (Lashmore-Davies 1975) which treats the pump and the excited waves on the same footing has been employed in our analysis to obtain the amplitude of the threshold electric field required for the onset of instability and the growth rate of the unstable Raman mode well above the threshold has been studied. A large magnetic field is found to enhance the threshold pump amplitude as well as the growth rate. It can also be seen from our analysis that a right-handed circularly polarized pump wave gives rise to a left-handed circularly polarized scattered wave as a result of the scattering from the electron plasma mode.

2. Theoretical formulation

Let us consider an infinite and isotropic semiconductor-plasma to which an uniform and static magnetic field is applied such that $\mathbf{B}_0 = B_0\,\mathbf{x}$. The plasma is subjected to a high frequency electromagnetic wave which is circularly polarized in a plane perpendicular to $\mathbf{B}_0$ and propagates in the x-direction. The scattering is completely characterized by the following wave vector and frequency selection rules:

$$\mathbf{k}_{T0} = \mathbf{k}_{T1} + \mathbf{k}_l, \tag{1}$$

and

$$\omega_{T0} = \omega_{T1} + \omega_l. \tag{2}$$

We have assumed the spatial uniformity and the perfect frequency matching and hence we take equations (1) and (2) to be satisfied exactly. Here $\mathbf{k}_{T0}$ ($\mathbf{k}_{T1}$) and ω_{T0} (ω_{T1}) are, respectively, the incident pump (scattered) wave vector and frequency, $\mathbf{k}_l$ and ω_l are the wave vector and frequency of the electron plasma wave that is excited in the semiconductor. The subscripts T_0, T_1 and l stand for the pump, scattered wave and Langmuir wave respectively. The electric field of the large amplitude circularly polarized pump wave is described by

$$E_{T0\pm} = E_{T0} \exp\left[i\left(k_{T0}\,x - \omega_{T0}\,t\right)\right].$$

The first step towards studying the parametric interaction of the two electromagnetic and an electron plasma wave is to obtain the normal modes for these waves. For this the starting equations are:

$$\frac{\partial v_\pm}{\partial t} + \nu v_\pm + \frac{e}{m} E_\pm \mp i\omega_c v_\pm = \mp \frac{ie}{m} v_x B_\pm - v_x \frac{\partial v_\pm}{\partial x}, \tag{3}$$

$$\frac{i\partial E_\pm}{\partial x} \pm \frac{\partial B_\pm}{\partial t} = 0, \tag{4}$$

$$\frac{\epsilon\partial E_\pm}{\partial t} \mp \frac{i\partial B_\pm}{\mu_0\,\partial x} - n_0\, e\, v_\pm = n_1\, e\, v_\pm, \tag{5}$$

$$\frac{\partial n_1}{\partial t} + n_0 \frac{\partial v_x}{\partial x} = -\frac{\partial}{\partial x}(n_1 v_x), \tag{6}$$

$$\frac{\partial v_x}{\partial t} + \nu v_x + \frac{eE_x}{m} + \frac{v_{\text{Th}}^2}{n_0}\frac{\partial n_1}{\partial x} = -\frac{e}{m}\left[\frac{i v_+ B_-}{2} - \frac{i v_- B_+}{2}\right] + \frac{v_{\text{Th}}^2\, n_1}{n_0}\frac{\partial n_1}{\partial x} - \frac{v_x\,\partial v_x}{\partial x} \tag{7}$$

and
$$\epsilon\frac{\partial E_x}{\partial t} - n_0 e v_x = n_1 e v_x. \tag{8}$$

Equations (3), (4) and (5) are used to obtain the normal modes of the pump wave and the scattered electromagnetic wave while (6), (7) and (8) yield the normal mode of Langmuir wave (electron plasma wave). Equations (3) and (7) are the momentum transfer equations while (4), (5) and (8) are the Maxwell's equations. All of them are written in the component form corresponding to the electromagnetic and Langmuir modes. $v_\pm$ and v_x are the perturbed fluid velocities. ν is the electron collision frequency and ω_c $(=|e|B_0/m)$ denotes the electron cyclotron frequency. $E_\pm$, E_x and $B_\pm$ are the perturbed electric field components and perturbed magnetic induction respectively. Equation (6) is the continuity equation. n_0 and n_1 denote the unperturbed and perturbed electron density respectively. $v_{\text{Th}}=(2K_BT_e/m)^{\frac{1}{2}}$ represents the thermal velocity of the electrons. The terms on the right side of (3) to (8) denote the nonlinear contributions. To obtain the normal mode equations we neglect the quadratic terms in the wave variables and take the linear combination of (3) to (5) and (6) to (8) such that

$$\frac{\partial}{\partial t}(v_{T\pm} \pm \alpha_T B_{T\pm} + \epsilon\beta_T E_{T\pm}) + \nu\, v_{T\pm} + \frac{eE_{T\pm}}{m} \mp i\omega_c\, v_{T\pm} - \alpha_T k_T E_{T\pm} \pm \frac{\beta_T k_T B_{T\pm}}{\mu_0} - \beta_T e n_0\, v_{T\pm} = 0, \tag{9}$$

and
$$\frac{\partial}{\partial t}(n_1 + \alpha_l v_x + \beta_l\, \epsilon E_x) + ik_l n_0 v_x + \alpha_l \nu v_x + \alpha_l\frac{eE_x}{m} + \frac{i\alpha_l v_{\text{Th}}^2 k_l n_1}{n_0} - \beta_l e n_0 v_x = 0, \tag{10}$$

where we have assumed $\partial/\partial_x = ik_{T,l}$. The normal mode equations are obtained by choosing $\alpha_{T,l}$ and $\beta_{T,l}$ such that equations (9) and (10) have the form

$$\frac{\partial a_{T\pm}}{\partial t} + i\omega_T a_{T\pm} + \gamma_T a_{T\pm} = 0, \tag{11a}$$

and $$\frac{\partial a_l}{\partial t} + i\omega_l a_l + \gamma_l a_l = 0. \tag{11b}$$

$a_{T\pm}$ and a_l are the normal mode amplitudes represented as

$$a_{T\pm} = v_{T\pm} \pm \alpha_T B_{T\pm} + \epsilon\beta_T E_{T\pm}, \tag{12a}$$

$$a_l = n_1 + \alpha_l v_x + \epsilon\beta_l E_x, \tag{12b}$$

where we assume

$$a_{T\pm, l} \sim \exp\left[i(k_{T,l}x - \omega_{T,l}t)\right] \exp\left[-\gamma_{T,l}t\right].$$

Here γ_T and γ_l represent the collisional damping rates. The subscript T denotes electromagnetic waves and subscripts $\pm$ represent the nature of polarization. From (10) and (11), the constants $\alpha_{T,l}$, $\beta_{T,l}$ and $\gamma_{T,l}$ are obtained as

$$\alpha_T = -\frac{k_T}{\mu_0 e n_0}\left[\frac{1 \pm \dfrac{\omega_c}{\omega_T} + \dfrac{i\nu}{\omega_T}\left\{1 - \dfrac{\omega_p^2\,\omega_T}{\{2\omega_T(\omega_T \pm \omega_c)^2 \mp \omega_p^2\,\omega_c\}}\right\}}{\left\{1 - \dfrac{i\nu\omega_p^2\,\omega_T}{\{2\omega_T(\omega_T \pm \omega_c)^2 \mp \omega_p^2\,\omega_c\}}\right\}}\right], \tag{13a}$$

$$\beta_T = -\frac{i\omega_T}{e n_0}\left[1 \pm \frac{\omega_c}{\omega_T} + \frac{i\nu}{\omega_T}\left\{1 - \frac{\omega_p^2\,\omega_T}{\{2\omega_T(\omega_T \pm \omega_c)^2 \mp \omega_p^2\omega_c\}}\right\}\right], \tag{13b}$$

$$\gamma_T = \frac{\nu\omega_p^2\,\omega_T}{2\omega_T(\omega_T \pm \omega_c)^2 \mp \omega_p^2\,\omega_c}, \tag{13c}$$

$$\beta_l = -\frac{ien_0}{\epsilon m v_{\mathrm{Th}}^2 k_l}, \tag{13d}$$

$$\alpha_l = \frac{n_0\omega_l}{v_{\mathrm{Th}}^2 k_l}\left(1 - \frac{i\nu}{\omega_l}\right), \tag{13e}$$

and $$\gamma_l = \nu/2, \tag{13f}$$

where we have used the dispersion relations

$$\omega_T^2 = k_T^2 c_1^2 + \frac{\omega_p^2\,\omega_T}{(\omega_T \pm \omega_c)} \text{ and } \omega_l^2 = \omega_p^2 + k_l^2 v_{\mathrm{Th}}^2.$$

C_1 is the velocity of light in the crystal. We take the same linear combinations of equations (3)–(5) and (6)–(8) as done for the normal modes, and retain only those nonlinear terms on the right side which have the space and time-dependence of the mode concerned. For simplicity of calculation, we have assumed that the incident pump wave is right-handed circularly polarized so that the scattered electromagnetic mode is a left-handed circularly polarized wave. Thus using (3)–(11), we obtain the coupled mode equations as

$$\frac{\partial a^*_{T1-}}{\partial t} + i\omega_{T1}\, a^*_{T1-} + \gamma_{T1}\, a^*_{T1-} = \beta^*_{T1-}\, en_1\, v^*_{T0^+} \mp \frac{iev_x\, B^*_{T0^+}}{m} - v_x \frac{\partial v^*_{T0^+}}{\partial x} \tag{14a}$$

and

$$\frac{\partial a_l}{\partial t} + i\omega_l a_l + \gamma_l a_l = \alpha_l \frac{e}{m}\frac{1}{2}(iv^*_{T1-}\;\; B_{T0^+}). \tag{14b}$$

These equations have been considered for a^+_T and a^+_l only, since these modes do not couple with a^-_l and a^-_T in view of the fact that the modes are of high frequencies and the coupling is weak. The equations for a^-_l and a^-_T are of the same form as (14) and do not give any additional information. Thus (14) covers both the cases of forward and backward propagating modes respectively. It is assumed that the mode amplitudes are small and the dependence of original variables ($v_{T\pm}$, n_1, etc.) on the mode amplitudes $a_{T\pm}$ and a_l is obtained using linearized forms of equations (3) to (8) and the definition of normal modes equation (12). The effect of collisions on the nonlinear terms is neglected since we have considered only weakly damped waves.

The normal mode amplitudes are taken as a product of a slowly varying amplitude and a rapidly varying phase as

$$a_{T\pm,\,l} = A_{T\pm,\,l}\,(t)\exp\,[i\,(k_{T,\,l}\,x - R_e\,\omega_{T,\,l}\,t)],$$

where the imaginary part of $\omega_{T,\,l}$ is contained in the mode amplitudes $A_{T,\,l}$. Using the above form in (14a, b) and dividing by the corresponding phase factors (so that $R_e\;\omega_{T,\,l}$ are completely eliminated, and what remains is the imaginary part denoted as ω) one gets the following dispersion relation

$$(-\,i\,\omega + \gamma_{T1})\,(-\,i\,\omega + \gamma_l) = C_{0l}\;C_{01}\,|\,A_{T0\,+}\,|^2, \tag{15}$$

where it is assumed that

$$A_{T0+} = \text{constant},\; |\,A_{T0+}\,| \gg |\,A_{T1-}\,|,\quad |\,A_{T0+}\,| \gg |\,A_l\,|. \tag{16}$$

In order to obtain the coupling terms C_{0l} and C_{01}, we substitute the values of the wave variable ($v,_{T0+}$ n_1, etc.) in (14) and we get

$$C_{0l} = \frac{v^2_{\text{Th}}\, k_l\, \omega^2_p}{2n_0\,\omega^2_l\,\{2\omega_{T0}\,(\omega_{T0} + \omega_c)^2 - \omega^2_p\,\omega_c\}}\,[\omega_{T0}(\omega_{T1}-\omega_c)\,k_l - \omega_l\,\omega_c\,k_{T0}]$$

and

$$C_{01} = \frac{\omega^4_p\, n_0\, \omega_l\, k_{T0}\, \omega_{T1}\,(\omega_{T0} + \omega_c)}{2v^2_{\text{Th}}\, k_l\,\{2\omega_{T0}\,(\omega_{T0} + \omega_c)^2 - \omega^2_p\,\omega_c\}\,\{2\omega_{T1}\,(\omega_{T1} - \omega_c)^2 + \omega^2_p\,\omega_c)\}}.$$

Equation (15) is solved analytically to obtain the threshold condition for growing Langmuir and scattered electromagnetic wave. The threshold electric field for the onset of instability is obtained by putting $\omega = 0$ in (15) as a result

$$(E_{\text{th}})_{B_0 \neq 0} = \frac{m\nu\omega_{T0}}{e}\left[\frac{2(\omega_{T0}+\omega_c)\,\omega_l}{k_{T0}\{\omega_{T0}(\omega_{T1}-\omega_c)\,k_l-\omega_c\,\omega_l\,k_{T0}\}}\right]^{1/2}. \quad (17)$$

To obtain the growth rate well above the threshold, we neglect the damping terms (γ_{T1}- and γ_l) in (15) and is solved for ω, then

$$(\omega)_{B_0 \neq 0} = \frac{e\,\omega_p\,E_{T0}}{2m\,\omega_{T0}}\left[\frac{\omega_{T1}\,k_{T0}\{\omega_{T_0}(\omega_{T1}-\omega_c)\,k_l-\omega_c\,\omega_l\,k_{T0}\}}{(\omega_{T0}+\omega_c)\,\omega_l\,\{2\omega_{T1}(\omega_{T1}-\omega_c)^2+\omega_p^2\,\omega_c\}}\right]^{1/2}. \quad (18)$$

In (17) and (18) we have assumed that

$$\omega_{T0,1} \sim \omega_l\,(\approx \omega_p) \sim \omega_c.$$

If we approximate

$$\omega_{T0} \approx \omega_{T1},\ k_{T0} \sim k_l \sim k$$

equations (17) and (18) give

$$(E_{\text{th}})_{B_0 \neq 0} = \frac{m\,\nu\,\omega_{T0}}{ek}\left[\frac{(\omega_{T0}+\omega_c)\,\omega_p}{\{\omega_{T0}(\omega_{T0}-\omega_c)-\omega_c\,\omega_p\}}\right]^{1/2}, \quad (19a)$$

and

$$(\omega)_{B_0 \neq 0} = \frac{ek\,E_{T0}}{2m}\left[\frac{\omega_p\,\{\omega_{T0}(\omega_{T0}-\omega_c)-\omega_c\,\omega_p\}}{\omega_{T0}(\omega_{T0}+\omega_c)\,\{2\omega_{T0}(\omega_{T0}-\omega_c)^2+\omega_p^2\,\omega_c\}}\right]^{1/2}. \quad (19b)$$

Equations (19a) and (19b) give the threshold pump amplitude required for the onset of instability and the growth rate for the case when the pressure term (v_{Th}^2 in (7)) has been included in the analysis.

For an isotropic plasma (*i.e.* $B_0 = 0$), (19a) and (19b) yield

$$(E_{\text{th}})_{B_0 = 0} = \frac{m\nu}{ek}\,[2\omega_{T_0}\,\omega_p]^{1/2}, \quad (20a)$$

and

$$(\omega)_{B_0 = 0} = \frac{ek\,E_{T0}}{2m\,\omega_{T0}}\,[\omega_p/2\omega_{T0}]^{1/2}, \quad (20b)$$

where it is assumed that $\omega_l \approx \omega_p$. For a magnetoactive plasma (*i.e.* $B_0 \neq 0$), we get from (19 a, b) and (20 a, b)

$$\frac{(E_{\text{th}})_{B_0 \neq 0}}{(E_{\text{th}})_{B_0 = 0}} = \left[\frac{\omega_{T0}(\omega_{T0}+\omega_c)}{\{\omega_{T_0}(\omega_{T0}-\omega_c)-\omega_c\,\omega_p\}}\right]^{1/2}, \quad (21a)$$

and
$$\frac{(\omega)_{B_0 \neq 0}}{(\omega)_{B_0 = 0}} = \omega_{T0} \left[\frac{2\{\omega_{T0}(\omega_{T0}-\omega_c)-\omega_c\,\omega_p\}}{(\omega_{T0}+\omega_c)\{2\omega_{T0}(\omega_{T0}-\omega_c)^2+\omega_p^2\,\omega_c\}} \right]^{1/2}. \tag{21b}$$

Neglecting the pressure term in (7) and following a similar procedure if the expressions for the threshold pump electric field and the growth rate of the unstable scattered electromagnetic and the Langmuir modes are obtained, it is found that they are the same as equations (19a) and (19b). Thus it can be concluded that the threshold conditions are insensitive to the electron thermal velocity v_{Th} in isotropic as well as magnetoactive plasmas.

3. Results and discussions

The analytical results obtained have been invoked to study the dependence of the threshold electric field and the growth rate on different parameters such as the pump wave frequency, carrier concentration, cyclotron frequency, thermal velocity, etc. It can be inferred from (19)–(21) that (i) the threshold electric field amplitude can be lowered by reducing the Langmuir wave frequency $\omega_l(\approx\omega_p)$. SRS from electron plasma waves in a semiconductor occurs only when ω_p is far away from the optical phonon frequency ω_{ph}. E_{th} can be reduced further by increasing the wave number k_l such that it satisfied the condition $k_l L \ll 1$ (where L is the electron mean free path) for the validity of hydrodynamic model of a semiconductor-plasma. (ii) Increasing longitudinal magnetostatic field increases the threshold pump amplitude as well as the growth rate. But at very high magnetic fields which yield $\omega_{T0}(\omega_{T0}-\omega_c) < \omega_c\omega_p$ in (21a) and (21b), the instability does not exist. (iii) The growth rate can be increased at higher values of wave number k and can be enhanced further by taking a pump wave with lower frequency such that it satisfies the frequency matching condition (equation (2)) and by using a semiconductor with smaller effective electronic mass and higher carrier concentrations. (iv) When the magnetic field is such that

$$0 < \frac{\omega_c}{\omega_{T0}} < 0.1$$

the threshold electric field amplitude and the growth rate differ very little from those at the vanishing magnetic field, (v) Since k_1 is always greater for the backscattering (*i.e.* $k_1 = k_{T0} + k_{T1}$), it has a much lower threshold than that for the forward scattering. (vi) Assuming $\omega_{T0,1} > \omega_p$ and using (21a) and (21b) one obtains

$$(E_{\mathrm{th}})_{B_0\neq 0}/(E_{\mathrm{th}})_{B_0=0} = [(\omega_{T0}+\omega_c)/(\omega_{T0}-\omega_c)]^{1/2} \tag{22a}$$

and
$$(\omega)_{B_0\neq 0}/(\omega)_{B_0=0} = [\omega_{T0}/(\omega_{T0}^2-\omega_c^2)^{1/2}] \tag{22b}$$

Comparison of our results (equations (22a) and (22b)) with Maraghechi and Willett (1979) yields

$$(E_{\mathrm{th}})^A_{B_0\neq 0}/(E_{\mathrm{th}})^M_{B_0\neq 0} = 2\sqrt{2}\,[(\omega_{T0}+\omega_c)/(\omega_{T0}-\omega_c)]^{3/2} \tag{23}$$

Equation (23) shows that the threshold electric field amplitude obtained in our case is $2\sqrt{2}$ times that obtained by Maraghechi and Willett (1979) for an isotropic semiconductor plasma. The difference arises due to the fact that we have considered right-handed polarized pump wave giving rise to a left-handed polarized scattered wave both travelling in the forward direction whereas Maraghechi and Willett (1979) have taken a left-handed polarized pump wave and only a backscattered wave. (vii) Further it can be seen from (22a, b) that at $\omega_c \simeq \omega_{T0}$ the threshold as well as growth rate become infinite. (viii) At $\omega_p \simeq \omega_{ph}$ electron plasma as well as the optical phonons contribute to the stimulated Raman scattering and then one cannot consider the case of SRS from electron plasma mode only.

A numerical estimation of the threshold electric field required for the onset of the SRS and of the growth rate of the unstable modes at an electric field larger than the threshold electric field has been made by using the analytical results (20) and (22) for isotropic and the magnetoactive plasmas. We take an n-type InSb crystal at 77°K and pulsed CO_2 laser of 10·6 μm as a high frequency electromagnetic pump. When $k = 5\times10^6$ m^{-1}, $\omega_p = 0{\cdot}2\times10^{13}$ sec^{-1} (*i.e.* $n_0 = 3{\cdot}18\times10^{20}$ m^{-3}) and (collision frequency related to momentum transfer which is assumed to be constant) $=2{\cdot}52\times10^{11}$ sec^{-1} (obtained from the value of electron mobility $\mu=5\times10^5$ cm^2 v^{-1} sec^{-1} at 77°K used by Gershenzon *et al* 1974), $(E_{\text{th}})_{B=0}$ is equal to $3{\cdot}4\times10^5$ Vm^{-1}. It can be applied to n-InSb without any damage. The damage threshold can be increased by reducing the pulse duration. The growth rate $(\omega)_{B_0=0}$ is equal to 10^8 sec^{-1} for an electric field of 10^6 Vm^{-1}. In a magnetoplasma it is found that when applied dc magnetic field is increased from 1-8 tesla $(E_{\text{th}})_{B_0\neq0} / (E_{\text{th}})_{B_0=0}$ increases from 1 to 1.85 and $(\omega)_{B_0\neq0} / (\omega)_{B_0=0}$ increases from 1·0 to 1·2.

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References

Drake J F, Kaw P K, Lee Y C, Schmidt G, Liu C S and Rosenbluth M N 1974 *Phys. Fluids* **17** 778
Foo E N and Tzoar N 1970 *Phys. Rev.* **B1** 2177
Fuchs V 1976 *Phys. Fluids* **19** 1554
Gershenzon E M, Kurilenko V A, Litvak-Gorskaya L B and Ravinovich R I 1974 *Sov. Phys. Semicond.* **7** 1005
Kunze H J, Funer L, Kronast B and Kegel W H 1964 *Phys. Lett.* **11** 42
Lashmore-Davies C N 1975 *Plasma Phys.* **17** 281
Makarov V P 1969 *Sov. Phys. JETP* **28** 336
Maraghechi B and Willett J E 1979 *Plasma Phys.* **21** 163
Patel C K N and Shaw E D 1970 *Phys. Rev. Lett.* **24** 451
Sodha M S, Sharma R P and Kaushik S C 1976 *J. Appl. Phys.* **47** 3518
Wolff P A 1970 *Phys. Rev.* **B1** 950
Wright G B, Kelley P L and Groves S H 1968 *Light scattering spectra of solids*, Proc. Int. Conf., New York (New York: Springer-Verlag) p. 35