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Stimulated Brillouin scattering in magnetized diffusive semiconductor plasmas

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Abstract. Using the hydrodynamical model and following the coupled mode approach, detailed analytical investigation of stimulated Brillouin scattering is performed in an electrostrictive semiconductor. The total induced current density including diffusion current density and the effective Brillouin susceptibility are obtained under off-resonant laser irradiation. The analysis deals with the qualitative behaviour of the Brillouin gain and transmitted intensity with respect to excess doping concentration and magnetic field. Efforts are directed towards optimizing the doping level and magnetic field to achieve maximum Brillouin gain at pump intensities far below the optical damage threshold level. It is found that by immersing a moderately doped semiconductor in a sufficiently strong magnetic field in transverse direction, one can achieve resonant enhancement of Brillouin gain provided the generated acoustic mode lies in the dispersionless regime.

PACS. 72.30.+q High-frequency effects; plasma effects – 77.65.Bn Piezoelectric and electrostrictive constants – 42.65.An Optical susceptibility, hyperpolarizability – 42.65.Es Stimulated Brillouin and Rayleigh scattering

1 Introduction

Stimulated scattering processes are nonlinear (NL) interactions in which an incident wave is converted to a frequency up or down shifted scattered wave. The difference in the photon energy between the incident and scattered wave is supplied by or taken up by the NL medium. Various types of scattering processes are possible, each involve different types of internal excitation of the medium. Stimulated Brillouin scattering (SBS) involves interactions with sound waves in solids, liquids, or gases or ion-acoustic waves in plasma [1–6].

SBS is known to be a valuable probe of acoustic phonons in gases, liquids and solids. The acoustic waves generated in solids due to SBS are amongst the most intense high-frequency sound waves and this may sometimes damage the materials [7]. SBS has recently been receiving considerable attention owing to its numerous applications in diverse areas ranging from optical phase conjugation (OPC), real-time holography, pulse compression to laser induced fusion [8–10]. For, OPC, backward SBS is preferred over other competitive techniques because it initiates at low threshold pump intensity, suffers negligible frequency shifts and offers high conversion efficiency [8]. In laser-induced fusion experiments SBS is of great concern because it can significantly redirect the pump energy away from the target and thus adversely affects the energy absorption. It is therefore desirable to minimize, or by any means control the SBS processes in these experiments.

Although SBS has been studied for more than three decades, the theoretical predictions and experimental measurements are far apart [9–12]. Several experiments, with short laser pulses of low intensity, suggest that SBS starts below the theoretically estimated threshold value, whereas other experiments using high intensity radiation reveal that SBS signal levels saturate at much lower values than their theoretically predicted values. Hence, more comprehensive efforts in SBS theory are needed.

In most of the investigations of NL interactions the nonlocal effects such as diffusion of the excitation density that is responsible for the NL refractive index change has normally been ignored. The diffusion of carriers is however expected to have strong influence on the nonlinearity of the medium, particularly in high mobility semiconductors viz., III-V compound semiconductors. Therefore, inclusion of carrier diffusion in theoretical studies of NL phenomenon seems to be important from both the fundamental and applied view points, and has thus attracted the attention of many groups recently [13–16]. Using a hydrodynamic model of semiconductor plasmas we present a study of the SBS phenomenon through the third order optical susceptibility, originating from the finite induced

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current densities and electrostrictive (ES) polarization, in transversely magnetized n-type semiconductors, in which diffusion of charge carriers is a compulsory phenomena. We have introduced the diffusion of charge carriers by considering the total current density as the sum of conduction and diffusion current densities.

The motivation for the present study stems from the fact that the diffusion of excess carriers can remarkably modify the nonlinearity of the medium. In the wake of high power lasers, such an investigation with high mobility semiconductor plasmas becomes even more important because it may lead to a better understanding of the scattering mechanisms in plasma media and thus may prove to be a step forward towards filling the gap between theory and experimental observations. The semiconductor is assumed to be immersed in a magnetic field which is expected to lower the SBS threshold and enhance the Brillouin gain appreciably.

2 Theoretical formulation

This section deals with the theoretical formulation of the NL optical susceptibility, and from there the steady state Brillouin gain for the Stokes component of the scattered electromagnetic wave in a Brillouin active medium. Here we consider a sample of n-type nearly centrosymmetric semiconductors, viz., n-InSb immersed in a uniform magnetostatic field $\overrightarrow{B_s}$ applied along z-axis. The semiconductor is assumed to be the source of a homogeneous and infinite plasma which is subjected to an externally driven large amplitude spatially uniform electromagnetic wave (pump wave), a high frequency laser or microwave propagating along x-axis. The electric field of the spatially uniform pump wave is described by $\overrightarrow{E_0} = E_0 \exp(-i\omega_0 t)$. We have chosen a centrosymmetric crystal so that the nonlinearities due to piezoelectricity and electro-optical effects can safely be ignored in comparison with those due to electrostriction.

In the hydrodynamic regime $k_a l \ll 1$ (k_a the acoustic wave number, and l the carrier mean free path), the basic equations considered for the analysis are:

$$\frac{\partial \overrightarrow{\vartheta_0}}{\partial t} + \nu \overrightarrow{\vartheta_0} = -\frac{e}{m}\left[\overrightarrow{E_0} + \left(\overrightarrow{\vartheta_0} \times \overrightarrow{B_s}\right)\right] = -\frac{e}{m}\overrightarrow{E_{\mathit{eff}}} \quad (1)$$

$$\frac{\partial \overrightarrow{\vartheta_1}}{\partial t} + \left(\overrightarrow{\vartheta_0}\frac{\partial}{\partial x}\right)\overrightarrow{\vartheta}_1 + \nu\overrightarrow{\vartheta_1} = -\frac{e}{m}\left[\overrightarrow{E_1} + \left(\overrightarrow{\vartheta_1} \times \overrightarrow{B_s}\right)\right] \quad (2)$$

$$\frac{\partial n_1}{\partial t} + \vartheta_0 \frac{\partial n_1}{\partial x} + n_0 \frac{\partial \vartheta_1}{\partial x} + D\frac{\partial^2 n_1}{\partial x^2} = 0 \quad (3)$$

$$\frac{\partial^2 u}{\partial t^2} - \frac{c}{\rho}\frac{\partial^2 u}{\partial x^2} + 2\Gamma_a \frac{\partial u}{\partial t} = \frac{\gamma}{2\rho}\frac{\partial}{\partial x}(\overrightarrow{E_{\mathit{eff}}} \cdot \overrightarrow{E_1^*}) \quad (4)$$

$$\overrightarrow{P_{es}} = -\gamma \overrightarrow{E_{\mathit{eff}}}\frac{\partial u^*}{\partial x} \quad (5)$$

$$\frac{\partial \overrightarrow{E_x}}{\partial x} = -\frac{n_1 e}{\varepsilon} + \frac{\gamma}{\varepsilon_1}\overrightarrow{E_0}\frac{\partial^2 u^*}{\partial x^2}. \quad (6)$$

Equations (1) and (2) represent the zeroth and first order oscillatory fluid velocities of an electron with effective mass 'm' and charge 'e' in which ν is the electron collision frequency. $\overrightarrow{E_{\mathit{eff}}}$ represents the effective electric field which includes the Lorentz force $\left(\overrightarrow{\vartheta_0} \times \overrightarrow{B_s}\right)$ in the presence of an external magnetic field $\overrightarrow{B_s}$. Equation (3) is the continuity equation including diffusion effects, where n_0, n_1 and D are the equilibrium and perturbed carrier densities and diffusion coefficient respectively. Equation (4) represents the lattice motion in the crystal where ρ is its mass density, u the lattice displacement, γ the electrostrictive coefficient, Γ_a the phenomenological damping parameter of acoustic mode and c is the elastic constant. Equation (5) reveals that the acoustic wave generated due to electrostrictive strain modulates the dielectric constant and gives rise to a NL induced polarization $\overrightarrow{P_{es}}$. At very high frequencies of the field, which are quite large as compared to the frequencies of the motion of electrons in the medium, the polarization is determined by neglecting the interactions of the electrons with one another and with nuclei of the atoms. Thus the electric displacement in the presence of an external magnetostatic field is simply given by $\overrightarrow{D} = \varepsilon\overrightarrow{E_{\mathit{eff}}}$ [17]. The space charge field E_x is determined by the Poisson equation (6), where ε_1 is the dielectric constant of the crystal. In the above equations, we have neglected the effect due to $\left(\overrightarrow{\vartheta_0} \times \overrightarrow{B_1}\right)$ by assuming that the acoustic wave is propagating along such a direction of the crystal so as to produce a longitudinal electric field.

The interaction of the pump with the electrostrictively generated acoustic wave produces an electron density perturbation, which in turn drives an electron plasma wave and induces current density in the Brillouin active medium. In a doped semiconductor, this density perturbation can be obtained by using a standard approach adopted by Dubey and Ghosh [18]. Differentiating equation (3) and using equations (1) and (6), we obtain

$$\frac{\partial^2 n_1}{\partial t^2} + \nu\frac{\partial n_1}{\partial t} + \nu D\frac{\partial^2 n_1}{\partial x^2} + \overline{\omega_p^2} n_1 + \frac{e k_1^2 n_0 \gamma u^* E_{\mathit{eff}}}{m\varepsilon_1} = i k_1 n_1 \overline{E} \quad (7)$$

where $\overline{E} = \frac{e}{m}\overrightarrow{E_{\mathit{eff}}}$, $\overline{\omega_p^2} = \left[\omega_p^2\left(\frac{\nu^2}{\nu^2+\omega_c^2}\right)\right]$, $\overrightarrow{E_{\mathit{eff}}} = \overrightarrow{E_0} + \left(\overrightarrow{\vartheta_0} \times \overrightarrow{B_s}\right)$.

Here $\omega_c \left[= \left(\frac{eB_s}{m}\right)\right]$ is the cyclotron frequency and $\omega_p \left[= \left(\frac{n_0 e^2}{m\varepsilon}\right)^{1/2}\right]$ is the plasma frequency of carriers in the medium. We neglect the Doppler shift under the assumption that $\omega_0 \gg \nu \gg k_0\vartheta_0$.

As per the method adopted by Dubey and Ghosh [18], the perturbed electron density (n_1) produced in the medium may be divided into two components which may be recognized as fast and slow components. The fast component (n_{1f}) corresponds to the first order Stokes component of scattered light and varies as $\exp\left[i\ (k_1 x - \omega_1 t)\right]$, whereas the slow component (n_{1s}) is associated with the acoustic wave and varies as $\exp\left[i\ (k_a x - \omega_a t)\right]$.

The process of SBS may also be described as the annihilation of a pump photon and simultaneous creation of one scattered photon and one induced photon. Hence, for these modes, the stimulated Brillouin process under consideration should satisfy the phase-matching conditions $\hbar\omega_0 = \hbar\omega_1 + \hbar\omega_a$ and $\hbar\vec{k_0} = \hbar\vec{k_1} + \hbar\vec{k_a}$ known as the energy and momentum conservation relations which determine the frequency shift and direction of propagation of the scattered light. By assuming a long interaction path for the interacting waves we consider only the resonant stokes component $\left(\omega_1 = \omega_0 - \omega_a,\ \vec{k_1} = \vec{k_0} - \vec{k_a}\right)$, and neglect the off-resonant higher-order components [19]. Moreover, for a spatially uniform pump, we assume that $\vec{k_1} = \vec{k_0} - \vec{k_a} \approx -\vec{k_a}$: $\vec{k_0}$ is zero under the dipole approximation.

We obtain the following coupled equations from equation (7) under the rotating wave approximation (RWA):

$$\frac{\partial^2 n_{1f}}{\partial t^2} + \nu\frac{\partial n_{1f}}{\partial t} + \nu D\frac{\partial^2 n_{1f}}{\partial x^2} + \overline{\omega_p^2} n_{1f} + \frac{ek_1^2 n_0\,\gamma\,u^* E_{\mathit{eff}}}{m\varepsilon_1} = -ik_1\overline{E}n_{1s}^* \tag{8a}$$

and

$$\frac{\partial^2 n_{1s}}{\partial t^2} + \nu\frac{\partial n_{1s}}{\partial t} + \nu D\frac{\partial^2 n_{1s}}{\partial x^2} + \overline{\omega_p^2} n_{1s} = -ik_1\overline{E}n_{1f}^*. \tag{8b}$$

From the above equations, it may be inferred that the generated acoustic wave and the Stokes mode couple to each other via the pump electric field in an ES medium. Hence, it is obvious that the presence of the pump field is of fundamental necessity for SBS to occur.

The slow component n_{1s} may be obtained from equations (4) and (8) as

$$n_{1s} = \frac{\varepsilon_0 k_1 k_a n_0 \gamma^2 E_{\mathit{eff}} E_1^* \,[A]^{-1}}{2\rho\varepsilon\left(\delta_a^2 - 2i\Gamma_a\omega_a\right)} \tag{9}$$

where, $A = \left[1 - \frac{\left(\delta_1^2 - i\omega_1\nu\right)\left(\delta_2^2 + i\omega_a\nu\right)}{k_1^2\,\overline{E}^2}\right]$, $\delta_a^2 = \omega_a^2 - k_a^2\vartheta_a^2$, $\delta_1^2 = \left(\overline{\omega_p^2} - \omega_1^2 - k^2\nu D\right)$ and $\delta_2^2 = \left(\overline{\omega_p^2} - \omega_a^2 - k^2\nu D\right)$.

It is evident from the above expression that n_{1s} strongly depends upon the magnitude of the pump intensity. The density perturbation thus produced affects the propagation characteristics of the generated waves.

The Stokes component $\left(\omega_1, \vec{k_1}\right)$ of the induced current density may be obtained from the standard relation

$$J_1(\omega_1) = n_0 e\vartheta_{1x} + e\vartheta_{0x} n_{1s}^*. \tag{10}$$

The preceding analysis under RWA yields

$$J_1(\omega_1) = -\frac{\omega_p^2\,\nu\,\varepsilon\,E_1}{\left(\nu^2 + \omega_c^2\right)} - \frac{\omega_p^2\,\varepsilon_0\,\gamma\,k_1\,k_a\,E_{\mathit{eff}}^*\,E_0\,E_1\left(\nu - i\omega_0\right)\,[A]^{-1}}{2\,\rho\left(\delta_a^2 - 2i\omega_a\Gamma_a\right)\left(\omega_c^2 - \omega_0^2\right)}. \tag{11}$$

The first term of the above expression represents the linear component of the induced current density. The second term represents the NL coupling amongst the three interacting waves via the total NL current density including the diffusion current.

The induced polarization may be expressed as the time integral of the induced current density. The polarization $P_{cd}(\omega_1)$ may therefore be obtained from equation (11) as

$$P_{cd}(\omega_1) = \frac{\omega_0^3\,k_1\,k_a\,\omega_p^2\,\varepsilon_0\,\gamma^2\,|E_0|^2\,E_1}{2\,\omega_1\,\rho\left(\omega_0^2 - \omega_c^2\right)^2\left(\delta_a^2 - 2i\omega_a\Gamma_a\right)}[A]^{-1}. \tag{12}$$

The origin of the SBS process lies in that component of $P_{cd}(\omega_1)$ which depends on $|E_0|^2 E_1$.

Hence, the threshold pump amplitude for the onset of SBS may be obtained by setting $P_{cd}(\omega_1) = 0$ in equation (12) as

$$E_{0th} = \frac{m\left(\omega_0^2 - \omega_c^2\right)}{e\,k_1\,\omega_0^2}\left|\left(\delta_1^2 - i\nu\omega_1\right)^{1/2}\left(\delta_2^2 + i\nu\omega_a\right)^{1/2}\right|. \tag{13}$$

Therefore, the interaction between the pump and the centrosymmetric crystal will be dominated by the SBS phenomena at a pump power level well above the threshold field E_{0th}.

Using the standard relation between induced polarization $P_{cd}(\omega_1)$ at frequency ω_1 and Brillouin susceptibility $(\chi_B)_{cd}$, one may obtain

$$(\chi_B)_{cd} = \frac{\omega_0^3\,k_1\,k_a\,\omega_p^2\,\gamma^2}{2\,\omega_1\,\rho\left(\omega_0^2 - \omega_c^2\right)^2\left(\delta_a^2 - 2i\omega_a\Gamma_a\right)}[A]^{-1}. \tag{14}$$

From equation (14) it can be inferred that the Brillouin susceptibility depends upon material parameters such as equilibrium carrier density, diffusion coefficient etc. It is also found that $(\chi_B)_{cd}$ depends on the magnitude of externally applied magnetic field $\vec{B_s}$ through the cyclotron frequency ω_c.

Besides this Brillouin susceptibility, the system should also possess an ES polarization, which arises due to the interaction of the pump with the acoustic wave generated in the medium. The scattering of the pump wave from the acoustic phonons affords a convenient means of controlling the frequency, intensity and direction of a scattered beam. This type of control makes a large number of applications possible involving the transmission, display and processing of informations. The ES polarization is obtained from equation (5) as

$$\vec{P_{es}} = \frac{k_1\,k_a\,\gamma^2\,\omega_0^4\,|E_0|^2\,\vec{E_1}}{2\,\rho\left(\delta_a^2 - 2i\omega_a\Gamma_a\right)\left(\omega_0^2 - \omega_c^2\right)^2}. \tag{15}$$

Using equation (15) one may obtain the Brillouin susceptibility due to ES polarization as

$$(\chi_B)_{es} = \frac{\varepsilon_0\,k_1\,k_a\,\gamma^2\,\omega_0^4}{2\,\rho\left(\delta_a^2 - 2i\omega_a\Gamma_a\right)\left(\omega_0^2 - \omega_c^2\right)^2}. \tag{16}$$

From equations (14) and (16) we obtain the effective Brillouin susceptibility using the relation

$$(\chi_B)_{eff} = (\chi_B)_{cd} + (\chi_B)_{es} \tag{17}$$

as

$$(\chi_B)_{eff} = \frac{\varepsilon_0\, k_1\, k_a\, \gamma^2\, \omega_0^4\, (\delta_a^2 + 2\, i\, \omega_a\, \Gamma_a)}{2\, \rho\, (\delta_a^4 + 4\, \omega_a^2\, \Gamma_a^2)\, (\omega_0^2 - \omega_c^2)^2} \times \left[1 + \left(\frac{\omega_p^2}{\omega_1\, \omega_0}\right)(A)^{-1}\right]. \tag{18}$$

Rationalization of equation (18) yields the imaginary part of Brillouin susceptibility as

$$(\chi_{Bi})_{eff} = \frac{\varepsilon_0\, k_1\, k_a\, \gamma^2\, \omega_0^4\, \omega_a\, \Gamma_a}{2\, \rho\, (\delta_a^4 + 4\, \omega_a^2\, \Gamma_a^2)\, (\omega_0^2 - \omega_c^2)^2} \times \left[1 + \left(\frac{\omega_p^2}{\omega_1\, \omega_0}\right)(A)^{-1}\right]. \tag{19}$$

In order to investigate the effective Brillouin gain constant, we use the following relation given in reference [20]:

$$\left[g\,(\omega_1)_{eff}\right] = -\frac{k}{2\,\varepsilon_1}\,(\chi_{Bi})_{eff}\,|E_0|^2\,. \tag{20}$$

Substituting the imaginary part of effective Brillouin susceptibility from equation (19) in equation (20) we obtain

$$\left[g\,(\omega_1)_{eff}\right] = -\frac{\varepsilon_0\, k_1\, k_a\, \gamma^2\, \omega_0^4\, \omega_a\, \Gamma_a}{4\, \rho\, (\delta_a^2 + 4\, \omega_a^2\, \Gamma_a^2)\, (\omega_0^2 - \omega_c^2)^2} \times \left[1 + \left(\frac{\omega_p^2}{\omega_1\, \omega_0}\right)(A)^{-1}\right]|E_0|^2\,. \tag{21}$$

Equation (21) can be used to compute the growth of Brillouin scattered mode in centrosymmetric diffusive crystals.

If the sample length is 10 to 10^2 orders greater than the pump wavelength, following Simoda [21] one can easily use the expression for effective induced polarization $(P_{nl})_{eff}$ $[= P_{cd} + P_{es}]$ deduced for an infinite medium, to express the transmitted electric field amplitude E_T in a crystal of cell length L,

$$E_T = -\frac{i\, k_1\, L}{\varepsilon}\left|(P_{nl})_{eff}\,(\omega_1)\right|, \tag{22}$$

which can be further written as

$$E_T = -\frac{i\, \varepsilon_0\, k_1^2\, k_a\, L\, \gamma^2\, \omega_0^4\, |E_0|^2\, E_1}{2\, \rho\, \varepsilon\, (\delta_a^2 - 2\, i\, \omega_a\, \Gamma_a)\, (\omega_0^2 - \omega_c^2)^2} \times \left[1 + \left(\frac{\omega_p^2}{\omega_1\, \omega_0}\right)(A)^{-1}\right]. \tag{23}$$

The above equation can be employed to determine the transmitted intensity (I_T) as

$$I_T = \frac{\eta\, \varepsilon_0^2\, C\, k_1^4\, k_a^2\, L^2\, \gamma^4\, \omega_0^8\, |E_0|^4\, E_1^2}{8\, \rho^2\, \varepsilon^2\, |R|^2\, (\omega_0^2 - \omega_c^2)^4} \times \left[1 + \left(\frac{\omega_p^2}{\omega_1\, \omega_0}\right)(A)^{-1}\right]^2 \tag{24}$$

where $R = \left(\delta_a^2 - 2\, i\, \omega_a\, \Gamma_a\right)$ and C is the velocity of light.

3 Results and discussion

By analyzing equation (13), it can be realized that the external magnetic field and wave number have strong influence on the threshold field E_{0th} required for the onset of useful Brillouin scattering in the crystal. E_{0th} decreases with increasing ω_c and k_1. This trend has also been observed experimentally by Muravjov and Shastin [22] in the investigation of stimulated FIR emission under a crossed field configuration. Considering the factor $\delta_1\left(= \sqrt{\left(\overline{\omega_p^2} - \omega_1^2 - k^2 \nu D\right)}\right)$, it may be concluded that the threshold field for the stimulated process increases with increasing carrier number density and decreasing diffusion coefficient.

A close look at equation (18) reveals that the effective Brillouin susceptibility is a sensitive function of carrier concentration via plasma frequency ω_p, magnetic field through cyclotron frequency ω_c and diffusion coefficient through the factor A. The cubic Brillouin susceptibility with carrier density 10^{24} m^{-3} purely due to diffusion current is found to be $\approx 8.5 \times 10^{-19}$ esu. At lower concentration this magnitude of $(\chi_B)_{eff}$ decreases by about five orders of magnitude and becomes potentially non-usable for the fabrication of cubic NL devices. The magnitude of the third-order susceptibility due to total current density (conduction as well as diffusion) agrees reasonably with the experimentally observed [22] and theoretically quoted values [23] using conduction current only.

A detailed numerical analysis of Brillouin gain and transmitted intensity is also made in an electrostrictive doped III-V semiconductor crystal at 77 K. The crystal is assumed to be subjected to a 10.6 μm nano second CO_2 laser. The material constants are taken as: $m = 0.015 m_0$ (m_0 being the free electron mass), $\rho = 5.8 \times 10^3$ kg m^{-3}, $\nu = 3 \times 10^{11}$ s^{-1}, $\omega_0 = 1.78 \times 10^{14}$ s^{-1}, $\omega_a = 10^{12}$ s^{-1}, $\Gamma_a = 2 \times 10^{10}$ s^{-1}, $\eta = 3.9$ and $\vartheta_a = 4.8 \times 10^3$ m s^{-1}.

We now focus our attention on the physical parameters that affect the Brillouin gain. It is found that the Brillouin gain increases with the input pump amplitude (Fig. 1). It is clear that higher pump intensity will cause greater gain. Figure 2 shows the variation of the gain constant as a function of the magnetic field (in terms of $\frac{\omega_c}{\omega_0}$). It is a unique feature displayed in this figure that finite gain is obtained only when ω_c is close to ω_0. This behaviour may be utilized for the construction of magnetic

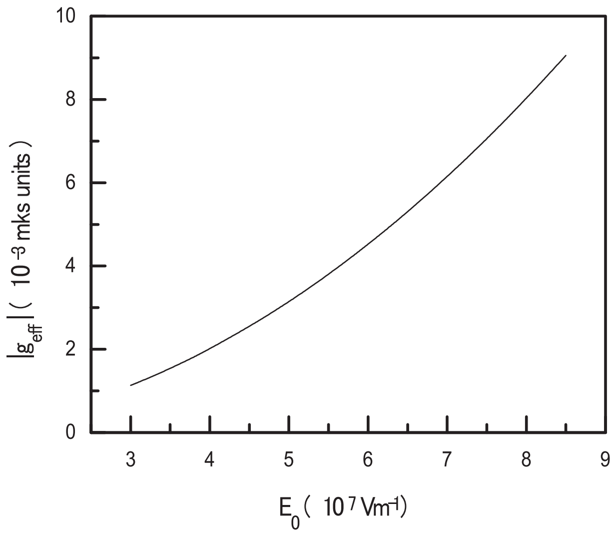

Fig. 1. Variation of gain with pump electric field.

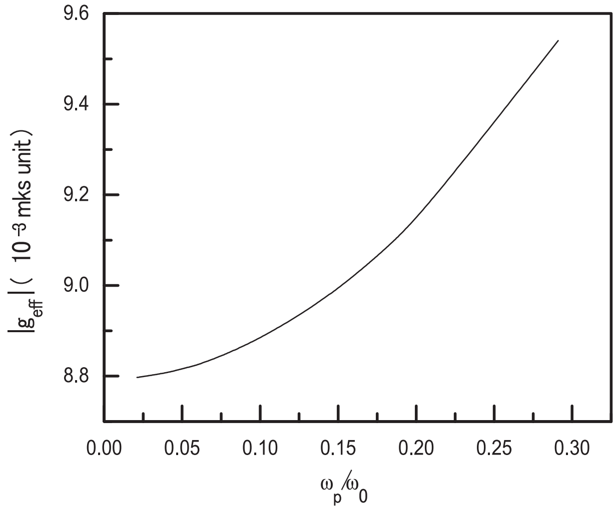

Fig. 3. Variation of gain with ω_p/ω_0.

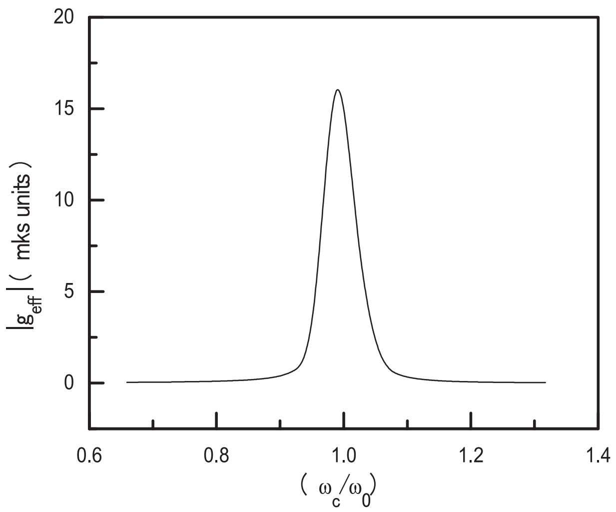

Fig. 2. Variation of gain with (ω_c/ω_0).

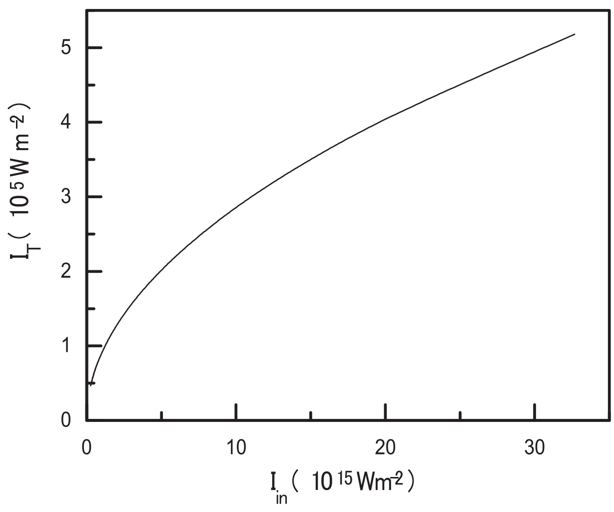

Fig. 4. Variation of transmitted intensity with input intensity.

switches. One can infer from this figure that the gain increases with the magnetic field up to $\omega_c \leqslant \omega_0$ and reaches its maximum value around $\omega_c \approx \omega_0$. The gain starts decreasing sharply in the region when $\omega_c \geqslant \omega_0$. This behaviour may be attributed to the fact that the gain coefficient is proportional to $\left(\omega_0^2 - \omega_c^2\right)^{-2}$ (Eq. (21)), and when $\omega_c > \omega_0$ magnetic absorption becomes vital which reduces the magnitude of the gain coefficient. The dependence of Brillouin gain on carrier number density (n_e) via plasma frequency ω_p is shown in Figure 3. It is clear from the figure that, as the carrier density increases, the gain increases rapidly. Figures 4–6 display the variation of transmitted intensity of the Brillouin scattered mode (I_T) with different input parameters when the cell length is taken as a hundred times larger than the input pump wave length. Figure 4 shows the nature of variation of transmitted intensity (I_T) with the input pump intensity. It is found that I_T increases rapidly with increasing pump intensity (I_{in}). Hence, in order to achieve maximum transmitted intensity of the Brillouin mode, it is always better to use a high intensity pump keeping damage threshold of the crystal in mind. The dependence of transmitted intensity on magnetic field (in terms of the ratio of cyclotron frequency with pump frequency) as shown in Figure 5 duplicates the dependency of Brillouin gain on $\frac{\omega_c}{\omega_0}$ as shown in Figure 2. Figure 6 shows that the transmitted intensity first increases with increasing doping level and then achieves a maximum value at about $\omega_p \approx 0.7\omega_0$. A little variation of doping level on the upper side ($\omega_p \geqslant 0.7\omega_0$) causes a drastic reduction in the magnitude of Brillouin gain until $\omega_p \approx 1.5\omega_0$ and then the gain marginally increases. This behaviour may be attributed to the factor within square brackets in equation (21). Hence, by adjusting the doping level one may maximize the Brillouin gain in the magnetized diffusive semiconductor very easily. This maximum gain is possible in moderately doped samples and hence carrier diffusion and magnetic field are favorable for the study of Brillouin gain in the crystal.

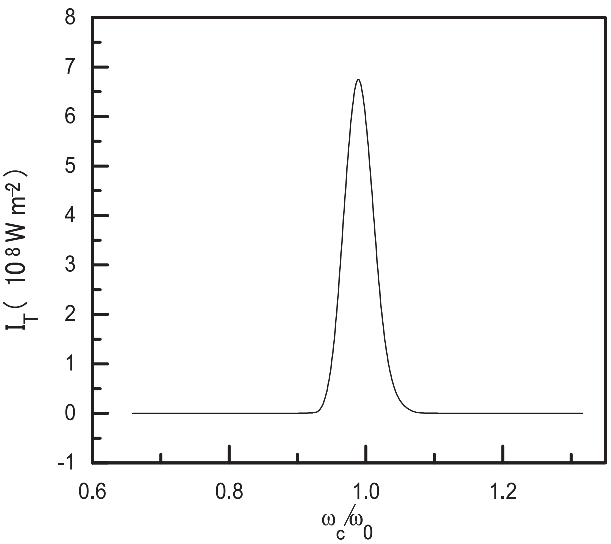

Fig. 5. Variation of transmitted intensity with ω_c/ω_0.

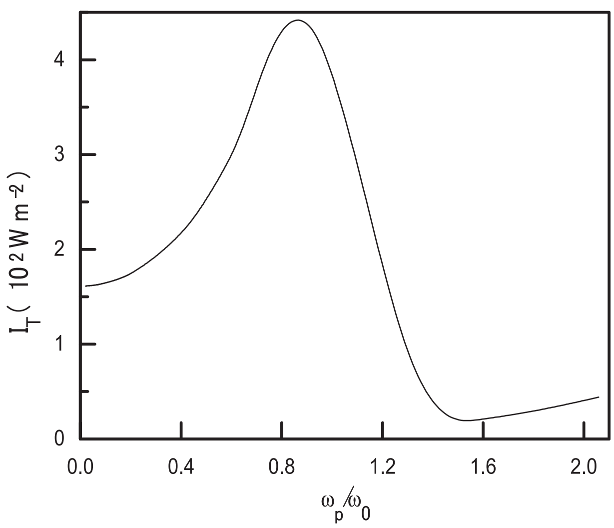

Fig. 6. Variation of transmitted intensity with ω_p/ω_0.

The large discrepancies between experimental and theoretical studies in solids may be attributed to the finite size of semiconductor plasmas as well as the finite values of drift velocities attainable in semiconductors, and the strong attenuating effects of scattering and Landau damping. The application of a magnetic field across the wave propagation direction tends to decrease the Landau damping effects. The above discussion reveals that large Brillouin gain and transmitted intensity can be easily achieved in moderately doped magnetized semiconductor plasmas. The present study thereby provides a model most appropriate for finite laboratory semiconductor plasmas, and an experimental study based on this work may provide new means for developing potentially useful Brillouin cells and for characterization and diagnosis of ES diffusive semiconductors.

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