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Original Research Article

Enhanced Brillouin gain in electrostrictive semiconductors using off-resonant laser excitation

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ABSTRACT
A detailed analytical study of stimulated Brillouin scattering (SBS) in an electrostrictive semiconductor is carried out using the hydrodynamic model in conjunction with the coupled-mode approach. The total induced current density, including contributions from diffusion, and the effective Brillouin susceptibility are derived for off-resonant laser excitation. The analysis examines the qualitative behavior of the Brillouin gain and transmitted intensity as functions of excess doping concentration and applied magnetic field. The study aims to optimize the doping level and magnetic field to achieve maximum Brillouin gain at pump intensities well below the optical damage threshold. It is observed that immersing a moderately doped semiconductor in a sufficiently strong transverse magnetic field can lead to resonant enhancement of the Brillouin gain, provided the generated acoustic mode lies within the dispersion-less regime.

1. Introduction

Stimulated scattering processes are nonlinear interactions in which an incident wave is transformed into a scattered wave with a frequency shifted either upward or downward. The energy difference between the incident and scattered photons is exchanged with the nonlinear medium. Different types of scattering processes can occur, each involving distinct internal excitations within the medium. In particular, stimulated Brillouin scattering (SBS) arises from interactions with acoustic waves in solids, liquids, or gases, or with ion-acoustic waves in plasmas [1–6].

Stimulated Brillouin scattering (SBS) serves as an important tool for probing acoustic phonons in gases, liquids, and solids. The acoustic waves generated in solids through SBS are among the most intense high-frequency sound waves, which can occasionally cause material damage [7]. Recently, SBS has garnered significant interest due to its wide range of applications, including optical phase conjugation (OPC), real-time holography, pulse compression, and laser-induced fusion [8–10]. For OPC, backward SBS is particularly advantageous because it operates at low threshold pump intensities, exhibits minimal frequency shifts, and provides high conversion efficiency [8]. In laser-induced fusion, however, SBS poses a challenge as it can deflect pump energy away from the target, thereby reducing energy absorption. Consequently, controlling or minimizing SBS is crucial in such experiments.

Despite over three decades of research on SBS, significant discrepancies remain between theoretical predictions and experimental observations [9–12]. Experiments with short, low-intensity laser pulses indicate that SBS can initiate below the theoretically predicted threshold, while studies employing high-intensity radiation show that SBS signals saturate at levels much lower than expected. These inconsistencies highlight the need for more comprehensive theoretical models of SBS.

In most studies of nonlinear interactions, nonlocal effects—such as the diffusion of excitation density, which contributes to changes in the nonlinear refractive index—are often neglected. However, carrier diffusion is expected to significantly influence the nonlinearity of the medium, particularly in high-mobility semiconductors, such as III-V compounds. Therefore, incorporating carrier diffusion into theoretical analyses of nonlinear phenomena is important from both fundamental and practical perspectives, and has recently attracted considerable attention [13–16]. Using a hydrodynamic model of semiconductor plasmas, we investigate the SBS phenomenon via the third-order optical susceptibility, arising from finite induced current densities and electrostrictive (ES) polarization, in transversely magnetized n-type semiconductors, where carrier diffusion is an inherent effect. The diffusion of charge carriers is included by expressing the total current density as the sum of conduction and diffusion contributions.

The motivation for the present study arises from the observation that the diffusion of excess carriers can significantly alter the nonlinearity of the medium. In the context of high-power laser interactions, investigating high-mobility semiconductor plasmas becomes particularly important, as it can provide deeper insights into scattering mechanisms in plasma media and help bridge the gap between theoretical predictions and experimental results. In this study, the semiconductor is assumed to be subjected to an external magnetic field, which is expected to reduce the SBS threshold and substantially enhance the Brillouin gain.

2. Theoretical formulation

This section presents the theoretical formulation of the nonlinear optical susceptibility and, based on it, the steady-state Brillouin gain for the Stokes component of the scattered electromagnetic wave in a Brillouin-active medium. We consider a sample of n-type, nearly centrosymmetric semiconductors, such as n-InSb, placed in a uniform magnetostatic field $\vec{B}_s$ applied along the *z*-axis. The semiconductor is treated as a source of a homogeneous, infinite plasma subjected to a large-amplitude, spatially uniform electromagnetic pump wave—either a high-frequency laser or microwave—propagating along the *x*-axis. The electric field of this spatially uniform pump wave is represented by $\vec{E}_0 = E_0 \exp(-i\omega_0 t)$. A centrosymmetric crystal is chosen so that nonlinearities arising from piezoelectric and electro-optical effects can be safely neglected relative to those resulting from electrostriction.

In the hydrodynamic regime, where $k_a l \ll 1$ (with k being the acoustic wave number and l the carrier mean free path), the fundamental equations used for the analysis are:

$$\frac{\partial \vec{v}_0}{\partial t} + \nu \vec{v}_0 = -\frac{e}{m}[\vec{E}_0 + (\vec{v}_0 \times \vec{B}_s)] = -\frac{e}{m}\vec{E}_{eff} \quad (1)$$

$$\frac{\partial \vec{v}_1}{\partial t} + \left(\vec{v}_0 \frac{\partial}{\partial x}\right)\vec{v}_1 + \nu \vec{v}_1 = -\frac{e}{m}[\vec{E}_1 + (\vec{v}_1 \times \vec{B}_s)] \quad (2)$$

$$\frac{\partial n_1}{\partial t} + v_0 \frac{\partial n_1}{\partial x} + n_0 \frac{\partial v_1}{\partial x} + D\frac{\partial^2 n_1}{\partial x^2} = 0 \quad (3)$$

$$\frac{\partial^2 u}{\partial t^2} - \frac{C}{\rho}\frac{\partial^2 u}{\partial x^2} + 2\Gamma_a \frac{\partial u}{\partial t} = \frac{\gamma}{2\rho}\frac{\partial}{\partial x}(\vec{E}_{eff} \cdot \vec{E}_1^*) \quad (4)$$

$$\vec{P}_{es} = -\gamma \vec{E}_{eff} \frac{\partial u^*}{\partial x} \quad (5)$$

$$\frac{\partial \vec{E}_x}{\partial x} = -\frac{n_1 e}{\varepsilon} + \frac{\gamma}{\varepsilon_1}\vec{E}_0 \frac{\partial^2 u^*}{\partial x^2}. \quad (6)$$

Equations (1) and (2) describe the zeroth- and first-order oscillatory fluid velocities of an electron with effective mass m and charge e, where ν is the electron collision frequency. $\vec{E}_{eff}$ represents the effective electric field, which includes the Lorentz force $(\vec{v}_0 \times \vec{B}_s)$ in the presence of an external magnetic field $\vec{B}_s$. Equation (3) is the continuity equation incorporating diffusion effects, with n_0 and n_1 denoting the equilibrium and perturbed carrier densities, respectively, and D the diffusion coefficient. Equation (4) governs the lattice motion in the crystal, where ρ is the mass density, u is the lattice displacement, γ is the electrostrictive coefficient, Γ_a is the phenomenological damping parameter of the acoustic mode, and C is the elastic constant. Equation (5) shows that the acoustic wave generated by electrostrictive strain modulates the dielectric constant, producing a nonlinear induced polarization $\vec{P}_{es}$. At very high field frequencies, which are much larger than the natural frequencies of electron motion in the medium, the polarization can be determined by neglecting electron–electron and electron–nucleus interactions. Consequently, the electric displacement in the presence of an external magnetostatic field is given by $\vec{D} = \varepsilon \vec{E}_{eff}$ [17]. The space charge field E_x is obtained from Poisson's equation (6), where ε_1 is the dielectric constant of the crystal. In these equations, the contribution from $(\vec{v}_0 \times \vec{B}_1)$ is neglected by assuming that the acoustic wave propagates along a crystal direction that produces a purely longitudinal electric field.

The interaction between the pump wave and the electrostrictively generated acoustic wave produces a perturbation in the electron density, which subsequently drives an electron plasma wave and induces a current density in the Brillouin-active medium. In a doped semiconductor, this density perturbation can be calculated using the approach outlined by Pravesh et al. [18]. By differentiating Equation (3) and applying Equations (1) and (6), we obtain:

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

where $\bar{E} = \frac{e}{m}\vec{E}_{eff}$, $\bar{\omega}_p^2 = \omega_p^2 \left(\frac{\nu^2}{\nu^2 + \omega_c^2}\right)$,

$\vec{E}_{eff} = \vec{E}_0 + (\vec{v}_0 \times \vec{B}_s)$.

Here, $\omega_c = \frac{eB_s}{m}$ denotes the cyclotron frequency,

and $\omega_p = \left(\frac{n_0 e^2}{m\varepsilon}\right)^{1/2}$ is the plasma frequency of carriers.

The Doppler shift is neglected under the assumption that $\omega_0 \gg \nu \gg k_0 v_0$.

Following the approach of Pravesh et al. [18], the perturbed electron density n_1 in the medium can be decomposed into two components: a fast and a slow component. The fast component n_{1f} corresponds to the first-order Stokes component of the scattered light and varies as $[i(k_1 x - \omega_1 t)]$, while the slow component n_{1s} is associated with the acoustic wave and varies as $\exp[i(k_a x - \omega_a t)]$.

The SBS process can also be described as the annihilation of a pump photon accompanied by the simultaneous creation of a scattered photon and an induced phonon. Accordingly, the stimulated Brillouin process must satisfy the phase-matching conditions $\hbar\omega_0 = \hbar\omega_1 + \hbar\omega_a$ and $\hbar\vec{k}_0 = \hbar\vec{k}_1 + \hbar\vec{k}_a$, which correspond to energy and momentum conservation and determine the frequency shift and propagation direction of the scattered light. Assuming a long interaction path for the waves, only the resonant Stokes component $(\omega_1 = \omega_0 - \omega_a, \vec{k}_1 = \vec{k}_0 - \vec{k}_a)$ is considered, while off-resonant higher-order components are neglected [19]. Furthermore, for a spatially uniform pump, $\vec{k}_1 = \vec{k}_0 - \vec{k}_a \approx -\vec{k}_a$ is assumed to be zero under the dipole approximation.

Applying the rotating wave approximation (RWA) to Equation (7), we obtain the following set of coupled equations:

$$\frac{\partial^2 n_{1f}}{\partial t^2} + \nu \frac{\partial n_{1f}}{\partial t} + \nu D \frac{\partial^2 n_{1f}}{\partial x^2} + \bar{\omega}_p^2 n_{1f} + \frac{e k_1^2 n_0 \gamma u^* E_{eff}}{m \varepsilon_1} = -i k_1 n_{1s}^* \bar{E} \quad (8a)$$

$$\frac{\partial^2 n_{1s}}{\partial t^2} + \nu \frac{\partial n_{1s}}{\partial t} + \nu D \frac{\partial^2 n_{1s}}{\partial x^2} + \bar{\omega}_p^2 n_{1s} = -ik_1 n_{1f}^* \bar{E} . \quad (8b)$$

From the above equations, it is evident that the generated acoustic wave and the Stokes mode are coupled through the pump electric field in an electrostrictive medium. Therefore, the presence of the pump field is fundamentally essential for the occurrence of SBS.

The slow component n_{1s} can be derived from Equations (4) and (8) as:

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

where, $A = \left[1 - \frac{(\delta_1^2 - i\nu\omega_1)(\delta_2^2 + i\nu\omega_a)}{k_1^2 \bar{E}^2} \right]$,

$\delta_a^2 = \omega_a^2 - k_a^2 v_a^2$,

$\delta_1^2 = (\bar{\omega}_p^2 - \omega_1^2 - k^2 \nu D)$, and

$\delta_2^2 = (\bar{\omega}_p^2 - \omega_a^2 - k^2 \nu D)$.

It is clear from the above expression that n_{1s} is strongly dependent on the pump intensity. The resulting density perturbation, in turn, influences the propagation characteristics of the generated waves.

The Stokes component $(\omega_1, \vec{k}_1)$ of the induced current density can be determined using the standard relation:

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

The above analysis, under the rotating wave approximation (RWA), gives:

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

In the above expression, the first term corresponds to the linear component of the induced current density, while the second term represents the nonlinear coupling among the three interacting waves through the total nonlinear current density, which includes the contribution from carrier diffusion.

The induced polarization can be expressed as the time integral of the induced current density. Consequently, the polarization $P_{cd}(\omega_1)$ can be obtained from Equation (11) as:

$$P_{cd}(\omega_1) = \frac{\omega_0^3 k_1 k_a \omega_p^2 \varepsilon_0 \gamma^2 |E_0|^2 E_1}{2\rho\omega_1(\delta_a^2 - 2i\Gamma_a \omega_a)(\omega_0^2 - \omega_c^2)^2} . \quad (12)$$

The SBS process originates from the component of P_{cd} (ω_1) that depends on $|E_0|^2 E_1$.

Therefore, the threshold pump amplitude for the onset of SBS can be determined by setting $P_{cd}(\omega_1) = 0$ in Equation (12), giving:

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

Consequently, at pump power levels well above the threshold field E_{0th}, the interaction between the pump and the centrosymmetric crystal is predominantly governed by the SBS phenomenon.

By employing the standard relation between the induced polarization P_{cd} (ω_1) at frequency ω_1 and the Brillouin susceptibility $(\chi_B)_{cd}$, one can write:

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

Equation (14) indicates that the Brillouin susceptibility depends on material parameters, including the equilibrium carrier density and the diffusion coefficient. It is also evident that $(\chi_B)_{cd}$ is influenced by the magnitude of the externally applied magnetic field $\vec{B}_s$ via the cyclotron frequency ω_c.

In addition to the Brillouin susceptibility, the system also exhibits electrostrictive (ES) polarization, which is generated through the interaction between the pump wave and the acoustic wave produced within the medium. The scattering of the pump wave by acoustic phonons provides an effective way to manipulate the frequency, intensity, and propagation direction of the scattered light. This capability enables numerous applications in areas such as information transmission, display technologies, and signal processing. The ES polarization can be derived from equation (5) as follows:

$$\vec{P}_{es} = \frac{k_1 k_a \gamma^2 \omega_0^4 |E_0|^2 E_1}{2\rho(\delta_a^2 - 2i\Gamma_a \omega_a)(\omega_0^2 - \omega_c^2)^2} . \quad (15)$$

From equation (15), the Brillouin susceptibility arising from the electrostrictive (ES) polarization can be determined as:

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

By combining equations (14) and (16), the effective Brillouin susceptibility can be determined using the following relation:

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

as

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

Simplifying equation (18) gives the imaginary component of the Brillouin susceptibility as:

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

To examine the effective Brillouin gain constant, we employ the relation provided in reference [20]:

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

By substituting the imaginary part of the effective Brillouin susceptibility from equation (19) into equation (20), we arrive at

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

Equation (21) allows for calculating the amplification of the Brillouin scattered mode in centrosymmetric diffusive crystals.

When the sample length exceeds the pump wavelength by 10 to 10^2 times, following Simoda [21], the expression for the effective induced polarization $(P_{nl})_{eff}$ [= P_{cd} + P_{es}], originally derived for an infinite medium, can be conveniently used to determine the transmited electric field amplitude E_T in a crystal of length L.

$$E_T = \frac{-ik_1L}{\varepsilon}\left|(P_{nl})_{eff}(\omega_1)\right|, \quad (22)$$

which can alternatively be expressed as

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

The above equation can be used to calculate the transmitted intensity I_T as

$$I_T = \frac{\eta\varepsilon_0^2 c k_1^4 k_a^2 L^2 \gamma^4 \omega_0^8 |E_0|^4 E_1^2}{8\rho^2\varepsilon^2\left|\delta_a^2 - 2i\Gamma_a\omega_a\right|^2(\omega_0^2-\omega_c^2)^4}\left[1+\left(\frac{\omega_p^2}{\omega_1\omega_0}\right)(A)^{-1}\right]^2 . \quad (24)$$

3. Results and discussion

An analysis of equation (13) reveals that both the external magnetic field and the wave number significantly affect the threshold field E_{0th} needed to initiate effective Brillouin scattering in the crystal. Specifically, E_{0th} decreases as ω_c and k_1 increase. Additionally, considering the relevant factor $\delta_1 = (\bar{\omega}_p^2 - \omega_1^2 - k^2\nu D)^{1/2}$, it can be inferred that the threshold field for the stimulated process rises with an increase in carrier density and a decrease in the diffusion coefficient.

A closer examination of equation (18) indicates that the effective Brillouin susceptibility is highly sensitive to the carrier concentration through the plasma frequency ω_p, the external magnetic field via the cyclotron frequency ω_c, and the diffusion coefficient through the factor A. For a carrier density of 10^{24} m^{-3}, the cubic Brillouin susceptibility arising solely from the diffusion current is approximately 8.5×10^{-19} esu. At lower carrier concentrations, this value decreases by roughly five orders of magnitude, rendering it potentially unsuitable for the development of cubic nonlinear devices. Furthermore, the third-order susceptibility due to the total current density (including both conduction and diffusion contributions) shows reasonable agreement with both experimental observations and previously reported theoretical values [22] calculated using only the conduction current.

A comprehensive numerical analysis of the Brillouin gain and transmitted intensity is also carried out for an electrostrictive, doped III-V semiconductor crystal at 77 K. The crystal is assumed to be illuminated by a 10.6 µm, nanosecond CO_2 laser. The material parameters used in the analysis are as follows: m = 0.015m_0 (m_0 being the free electron mass), ρ = 5.8×10^3 kg m^{-3}, ν = 3×10^{11} s^{-1}, ω_0 = 1.78×10^{14} s^{-1}, ω_a = 10^{12} s^{-1}, Γ_a = 2×10^{10} s^{-1}, η = 3.9 and v_a = 4.8×10^3 m s^{-1}.

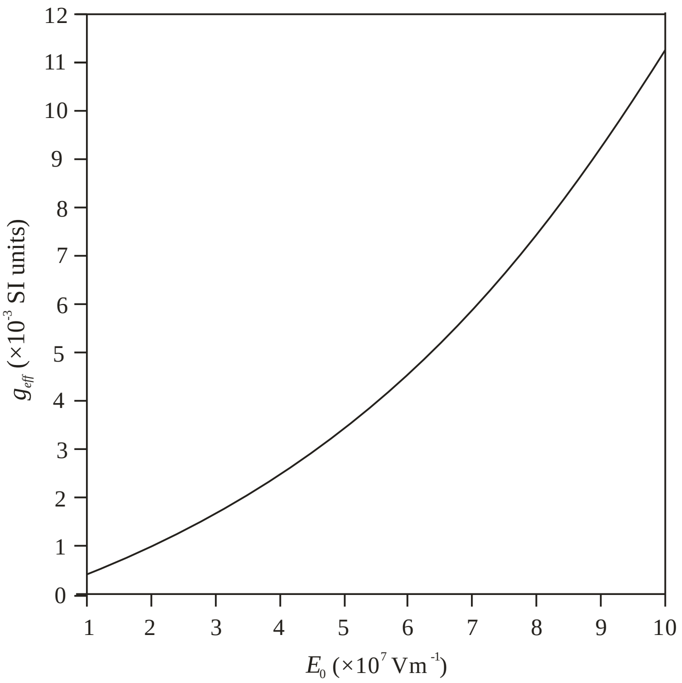

Figure 1: Dependence of Brillouin gain on the pump electric field.

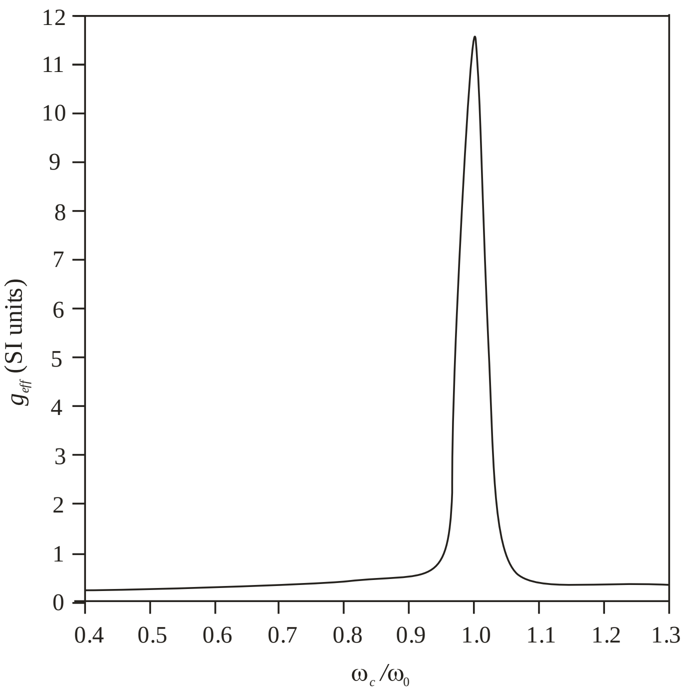

Figure 2: Brillouin gain as a function of the ratio ω_c/ω_0 .

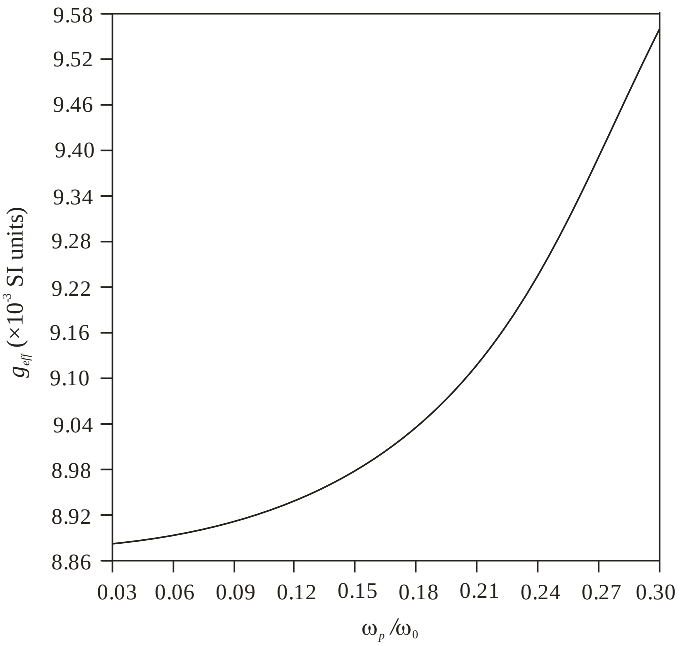

Figure 3: Brillouin gain as a function of the ratio ω_p/ω_0 .

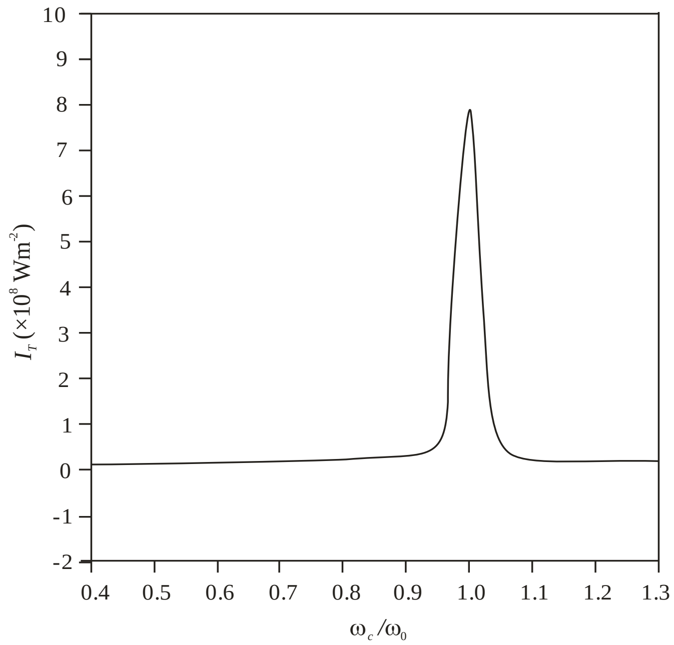

Figure 5: Transmitted intensity as a function of the ratio ω_c/ω_0 .

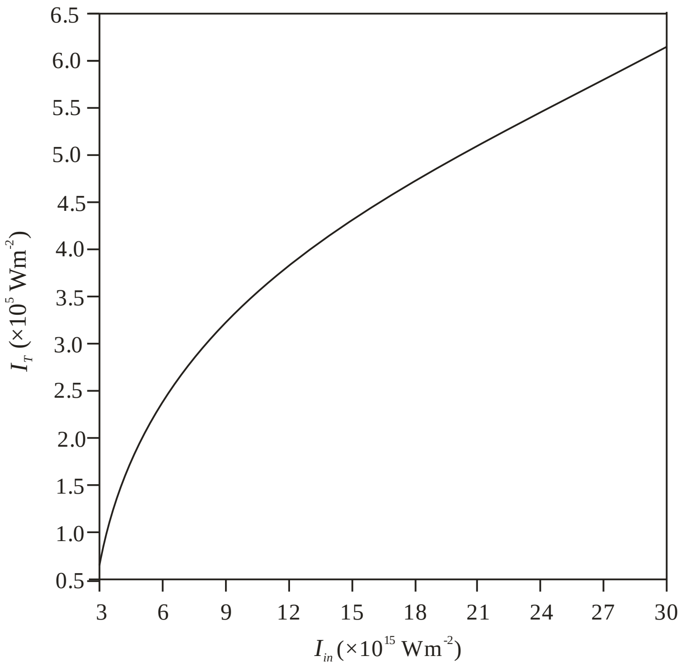

Figure 4: Transmitted intensity of the Brillouin scattered mode versus input pump intensity.

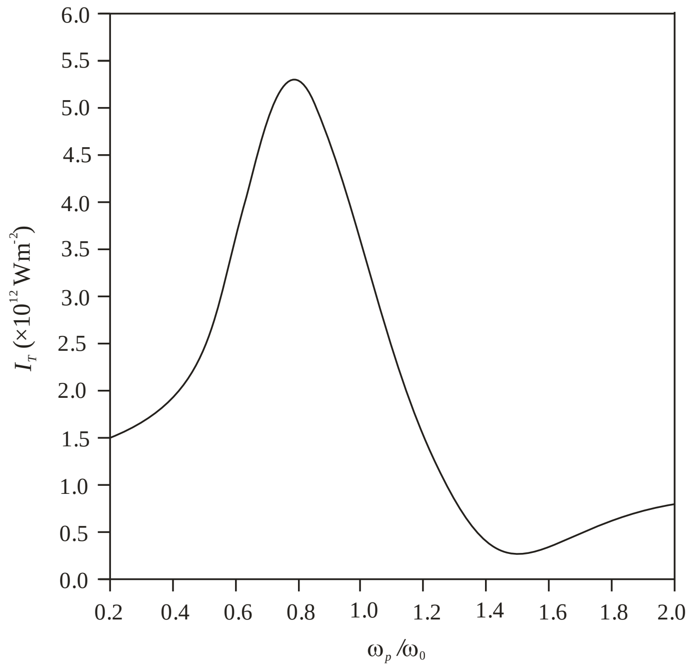

Figure 6: Dependence of transmitted intensity on the plasma frequency ratio ω_p/ω_0.

We now turn our attention to the physical parameters influencing the Brillouin gain. It is observed that the Brillouin gain increases with the input pump amplitude, as illustrated in Figure 1. This indicates that higher pump intensity results in greater gain. Figure 2 depicts the variation of the gain constant with the magnetic field, expressed in terms of ω_c/ω_0. A notable feature in this figure is that a finite gain is obtained only when ω_c is close to ω_0, a property that can be exploited in the design of magnetic switches. The gain rises with the magnetic field for $\omega_c < \omega_0$, reaching a maximum around $\omega_c \approx \omega_0$, and then decreases sharply when $\omega_c > \omega_0$. This behavior arises because the gain coefficient is proportional to the factor $(\omega_0^2 - \omega_c^2)^{-2}$ in equation (21), and magnetic absorption becomes significant when $\omega_c > \omega_0$, thereby reducing the gain.

The dependence of Brillouin gain on carrier density (n_e) via the plasma frequency ω_p is shown in Figure 3, indicating a rapid increase in gain with increasing carrier concentration. Figures 4–6 illustrate the transmitted intensity (I_T) of the Brillouin scattered mode under varying input parameters, assuming a crystal length 100 times the pump wavelength. Figure 4 shows that I_T rises sharply with increasing input pump intensity (I_{in}), suggesting that using a higher pump intensity enhances transmitted intensity, provided the crystal's damage threshold is not exceeded.

The dependence of transmitted intensity on the magnetic field, expressed as the ratio ω_c/ω_0 (Figure 5), mirrors the behavior of Brillouin gain shown in Figure 2. Figure 6 demonstrates that I_T initially increases with doping level, reaching a maximum at $\omega_p \approx 0.7\ \omega_0$. Further increases in doping beyond this point ($\omega_p > 0.7\omega_0$) lead to a sharp reduction in gain until $\omega_p \approx 1.5\omega_0$, after which the gain slightly recovers. This trend is associated with the factor enclosed in square brackets in equation (21). Therefore, careful adjustment of doping levels allows optimization of Brillouin gain in magnetized diffusive semiconductors. Maximum gain is achieved in moderately doped samples, highlighting the favorable roles of carrier diffusion and magnetic field in enhancing Brillouin scattering.

The significant differences observed between experimental and theoretical results in solids can be attributed to factors such as the finite size of semiconductor plasmas, the limited drift velocities achievable in semiconductors, and the pronounced attenuation caused by scattering and Landau damping. Applying a magnetic field perpendicular to the wave propagation direction helps to reduce the effects of Landau damping. The discussion above indicates that substantial Brillouin gain and transmitted intensity can be readily obtained in moderately doped, magnetized semiconductor plasmas. This study thus presents a model well-suited for finite laboratory semiconductor plasmas, and experimental investigations based on this framework could offer new avenues for the development of efficient Brillouin cells, as well as for the characterization and diagnostic analysis of electrostrictive diffusive semiconductors.

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Authors' contributions

The author read and approved the final manuscript.

Conflicts of interest

The author declares no conflict of interest.

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Data availability

No new data were created.

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