International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

Time Reversal Reflection in Diffusion Driven Semiconductor Plasmas: Effects of DC Magnetic Field and Doping Concentration

Sandeep
Department of Physics, Baba Mastnath University, Asthal Bohar, Rohtak, India
Sunita Dahiya
Department of Physics, Baba Mastnath University, Asthal Bohar, Rohtak, India
Navneet Singh
Department of Physics, A.I.J.H.M. College, Rohtak, India
Manjeet Singh
Department of Physics, Amity University, Noida, India

ABSTRACT

Time reversal reflection via stimulated Brillouin scattering (TRR-SBS) in magnetized diffusion driven semiconductors under the off-resonant transition regime has been investigated theoretically. The reflectivity of the image radiation is dependent upon the Brillouin susceptibility and can be significantly enhanced through n-type doping of the crystal and the simultaneous application of magnetic field. Moreover, the threshold of the pump intensity required for the occurrence of SBS in the crystal with finite optical attenuation can be considerably diminished through a suitable choice of the excess carrier concentration and the magnetic field. Numerical estimates made for n-InSb crystal at 77 K duly irradiated by nanosecond pulsed 10.6 μm CO_2 laser show that high TRR-SBS reflectivity (70%) can be achieved at pump intensities below the optical damage threshold if the crystal is used as an optical waveguide with relatively large interaction length (L ~5 mm) which proves its potential in practical applications such as fabrication of time reversal reflection mirrors.

Keywords

Time reversal reflection, stimulated Brillouin scattering, semiconductors, magnetic field, doping concentration

INTRODUCTION

Time reversal reflection (TRR) is a process that involves the use of nonlinear optical effects and precisely reverses the direction of propagation of each plane wave in an arbitrary beam of light, thereby causing the return beam to exactly retrace the path of the incident beam. The process is also known as wavefront reversal or optical phase conjugation. TRR has diverse potential applications in various areas ranging from adaptive optics, laser induced fusion, image transmission, pulse compression, image processing and real time holography [1-5]. The three major classes of nonlinear optical interactions used for TRR are optical parametric mixing, stimulated scattering processes and hybrid schemes [6]. The most commonly used techniques involving these interactions are: degenerate four-wave mixing (DFWM), stimulated Brillouin scattering (SBS) and Brillouin-enhanced four-wave mixing (BEFWM). TRR-SBS has been shown to have distinct advantages over TRR via other stimulated scattering processes owing to the facts that: (i) it is nearly a steady-state process; (ii) requires relatively low excitation intensity, and (iii) it suffers a negligibly frequency-shift [7].

The application of TRR-SBS can be widened significantly through the introduction of newer and newer configurations that enable one to achieve TRR-SBS reflectivity of the order of unity even at the reasonably low pump power available from a pulsed laser source with pulse duration larger than the acoustic phonon lifetime. The TRR-SBS reflectivity being directly dependent upon third-order (Brillouin) susceptibility $¡_B$,

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

an enhancement in reflectivity is possible even for not-too-high excitation intensity if one can achieve larger $¡_B$ in the medium. Recently, one of the present authors [8] have shown keen interest in the application of an external magnetic field to enhance remarkably the third-order optical susceptibility of the crystal. In the present paper, the authors have attempted to study the important phenomenon of time reversal reflection via stimulated Brillouin scattering in magnetized diffusion driven semiconductors duly irradiated by an intense uniform pump wave. The physical origin of the phenomenon lies in $¡_B$ arising due to the induced current density and acoustic perturbations internally generated due to electrortrictive property of the medium. The chief utility of the analysis is the fabrication of time reversal reflection mirrors with extremely high reflectivity.

TIME REVERSAL REFLECTIVITY

We consider the irradiation of the semiconductor crystal by a slightly off-resonant laser pump. The pump laser mode undergoes stimulated scattering processes via its interactions with phonons in the medium. Here, it is assumed that the scattering is due only to the acoustic phonons such that the interaction yields the scattered Brillouin mode. The propagation of these modes through the medium can be represented by the generalized equation [9]:

$$\frac{\partial^2 E(r_\perp,x,t)}{\partial x^2}+\nabla_t^2 E(r_\perp,x,t)+\frac{\mu_L^2}{c^2}E(r_\perp,x,t)+«_0\mu_L^2 P(r_\perp,x,t)=0 \qquad (1)$$

where $\nabla_t^2=(\partial^2/\partial y^2)+(\partial^2/\partial z^2)$; $E(r_\perp,x,t)$ is the electric field amplitude which possess both longitudinal and transverse components. The fields are assumed to vary as $\exp[i(\mu_L t-k_L x)]$, μ_L and k_L are the angular frequency and wave number of the electromagnetic mode, respectively; $«_0$ is the permeability of free space. P is total induced polarization in the crystal. The total optical susceptibility of the crystal is given by

$$P(r_\perp,x,t)=£_0 ¡E=£_0[¡^{(1)}+¡^{(3)}|E(r_\perp,x)|^2+...]E, \qquad (2)$$

where the crystal is assumed to be centrosymmetric such that the even-order nonlinear optical susceptibility components like $¡^{(2)}$, $¡^{(4)}$, etc. are zero. In bulk semiconductors, $¡^{(1)}$ does not play significant role in the improvement of conversion efficiency [10]. Nevertheless, it affects the spatial propagation characteristic of the conjugate beam. In the present analytical investigation of TRR-SBS, we have retained $¡$ comprising both linear as well as third-order components, i.e., $¡=¡^{(1)}+¡^{(3)}\left|E(r_\perp,x)\right|^2$.

The second term on the right hand side of Eq. (2) can be expressed in terms of a transmission function [10] proportional to $(E_0+E_s)(E_0^*+E_s^*)$, where E_0 and E_s are the electric fields associated with the pump (μ_0,k_0) and the scattered Stokes mode (μ_s,k_s), respectively. Consequently, the third-order component of the induced polarization is obtained as:

$$P^{(3)}=£_0¡^{(3)}[|E_0|^2+|E_s|^2+E_0^*E_s+E_0E_s^*]E \qquad (3)$$

We have assumed the pump wave (μ_0,k_0) to be propagated along the *–x* direction while the Stokes mode of the scattered electromagnetic wave is assumed to propagate along *+x* direction. Thus, one may replace k_L by $-k_0$ and k_s for the pump and the scattered mode, respectively. Making use of Eqs. (1) - (3), we can obtain the corresponding electromagnetic wave equations under the slowly varying envelope approximation as:

$$\frac{\partial E_0}{\partial x}-\frac{i}{2k_0}\nabla_t^2 E_0=\sim_{I0}E_0-i\sim_{Ir0}E_0-\frac{i\mu_0^2}{2k_0c^2}¡_0^{(3)}|E_s|^2E_0 \qquad (4a)$$

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

and

$$\frac{\partial E_s}{\partial x} - \frac{i}{2k_s}\nabla_t^2 E_s = -\sim_{IS} E_S + i\sim_{Irs} E_s + \frac{i\mu_s^2}{2k_s c^2}\,\mathrm{i}_s^{(3)}\left|E_0\right|^2 E_s \qquad (4b)$$

where $\sim_{I0}(\sim_{Is})$ is the intensity dependent absorption coefficient of the crystal at the pump frequency μ_0 (Stokes frequency μ_s) and given by

$$\sim_{I0,s} = \frac{\mu_{0,s}}{2c}\left[\mathrm{i}_{I0,s}^{(1)} + \mathrm{i}_{I0,s}^{(3)}\left|E_{s,0}\right|^2\right] \qquad (5a)$$

The parameters $\sim_{Ir0}$ and $\sim_{Irs}$ in Eq. (4) account for the dispersive properties of the material and have their origin in the real part of the optical susceptibility defined as:

$$\sim_{Ir0,s} = \frac{\mu_{0,s}}{2c}\left[\mathrm{i}_{r0,s}^{(1)} + \mathrm{i}_{r0,s}^{(3)}\left|E_{s,0}\right|^2\right]. \qquad (5b)$$

Moreover, $\mathrm{i}_{r0,s}^{(3)}$ in Eq. (4) represents the complex third-order susceptibility at frequency $\mu_{0,s}$. The phase matching conditions for TRR-SBS, as considered in the present study, are given by $\mu_0 = \mu_s + \mu_a$ and $\vec{k}_0 = -\vec{k}_s + \vec{k}_a$. For time reversal reflection Stokes mode, one should have $k_0 = -k_s$ such that $k_a = 2k_0$. In the forthcoming discussions, we have assumed $|k_0| = |k_s| = k$ such that $|k_a| = 2k$ and for considerably low acoustic frequency [i.e., $\mu_0 >> \mu_a$, one may take $\mu_0 \sim \mu_a = \mu$ (say)]. These assumptions enable one to take $\sim_{I0} = \sim_{Is} = \sim_I$ (say) and $\mathrm{i}_0^{(3)} \sim \mathrm{i}_s^{(3)} = \mathrm{i}_B$ with i_B being the Brillouin susceptibility. For stimulated Brillouin scattering, i_B is usually an imaginary quantity for dispersion-less acoustic-wave propagation [11, 12] and hence we take $\mathrm{i}_B = -i|\mathrm{i}_B|$. For a single pump and Stokes mode, we have followed the well-known single-mode formalism such that the time reversal reflection Stokes mode is related to the pump via relation

$$E_s(r_\perp, x) = \quad (x)E_0^*(r_\perp, x), \qquad (6)$$

where $\quad(x)$ is a measure of conjugacy with $\left|\quad(x)\right|^2$ being defined as TRR-SBS reflectivity.

Under one-dimensional configuration, the electric field associated with the pump and the scattered mode can be obtained from Eq. (4) as:

$$\frac{\partial E_0}{\partial x} = \sim_I E_0 - i\sim_{Ir} E_0 - \frac{i\mu^2}{2kc^2}\,\mathrm{i}_B\left|E_s\right|^2 E_0 \qquad (7a)$$

and

$$\frac{\partial E_s}{\partial x} = -\sim_I E_S + i\sim_{Ir} E_s + \frac{i\mu^2}{2kc^2}\,\mathrm{i}_B\left|E_0\right|^2 E_s \qquad (7b)$$

respectively, taking $\mu_0 \sim \mu_s \sim \mu = kc$. Since $|E_s|^2$ is a generated field, we may safely assume $\sim_I >> \mu^2\,\mathrm{i}_B\left|E_s(0)\right|^2 / 2kc^2$ in Eq. (7.7a). Consequently, we get a solution of equation (7a) as:

$$E_0(x) = \left[E_0(L)\exp\{-\sim_I(L-x)\}\right]\exp\{i\sim_{Ir}(L-x)\} \qquad (8)$$

where $E_0(L)$ is the pump electric field at the entrance window (i.e., at $x = L$). From Eqs. (7b) and (8), we find the electric field associated with the Stokes mode of the backscattered electromagnetic wave as:

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

$$E_s(x) = [E_s(0)\exp[-\{\sim_I x + \copyright\{1-\exp(2\sim_I x)\}\}/2\sim_I]\exp[i\sim_{Ir} x] \quad (9a)$$

where

$$\copyright = \frac{\mu^2 \mathfrak{i}_B}{2kc^2}|E_0(L)|^2 \exp(-2\sim_I L). \quad (9b)$$

In eq. (9a), $E_s(0)$ is the electric field amplitude at the exit window $x=0$ and can be defined as the spontaneous noise field. Now, Eq. (9a) can be rewritten in the form

$$E_s(x) = E_s'(0)[\cos(\sim_{Ir} L) + i\sin(\sim_{Ir} L)]\exp\{-i\sim_{Ir}(L-x)\} \quad (10a)$$

where

$$E_s'(0) = E_s(0)\exp[-\{\sim_I x + \copyright\{1-\exp(2\sim_I x)\}\}/2\sim_I] \quad (10b)$$

Eq. (10a) manifests the occurrence of the phenomenon of time reversal reflection in the material through the dependence of the scattered mode on the phase factor $\exp\{-i\sim_{Ir}(L-x)\}$ which is backscattered. Moreover, $\sim_{Ir}$ which depends upon the real part of $\mathfrak{i}_B$ will not affect either the pump or the scattered mode. As a consequence $|\ (x)|^2$ remains unaltered irrespective of the absolute phase difference between E_0 and E_s. Eq. (10b) can also be employed in studying the threshold nature of the TRR-SBS. The backscattered Brillouin mode $E_s(x)$ possesses a gain constant given by $\exp[-\{\sim_I x + \copyright\{1-\exp(2\sim_I x)\}\}/2\sim_I]$.

For finite gain, the condition

$$\left[\sim_I x + \frac{\copyright(1-e^{2\sim_I x})}{2\sim_I}\right] < 0. \quad (11a)$$

is to be achieved in the material medium.

For a bulk semiconducting crystal of mm thickness irradiated by a slightly off-resonant laser with photon energy less than the crystal band-gap energy or an optical fiber with very low-loss, one may take $2\sim_I x < 1$ such that the threshold condition Brillouin gain becomes

$$\copyright = \sim_I. \quad (11b)$$

The present work is restricted to the study of TRR-SBS in the semiconducting crystals irradiated by off-resonant (below the band-edge) laser sources. In this regime the third-order optical susceptibility is appreciably smaller than in the near-resonant case [13]. This enables one to treat $\sim_I$ and $\sim_{Ir}$ as the real and imaginary components of the background absorption coefficient $\sim$ and $\sim_r$, respectively, at frequency μ. Using Eqs. (9b) and (11) for $\sim L < 1$ at the entrance window $x = L$, the threshold value of excitation intensity is obtained as:

$$I_{0,th} = \frac{\text{¦}\pounds_0 c^3 k\sim}{\mu^2 \mathfrak{i}_B} \quad (12)$$

where ¦ being the background refractive index of the crystal, $\pounds_0$ is the absolute permittivity, $\sim$ is the background absorption coefficient at frequency μ, and $I_{0,th} = \frac{1}{2}\text{¦}\pounds_0 c|E_{0,th}|^2$. Again equations (8) and (10) give the TRR-SBS reflectivity as:

$$|\ (x)|^2 = \left[\frac{|E_s(0)|}{|E_0(L)|}\right]^2 \exp\left[2\left\{\sim(L-2x) - \frac{\copyright(e^{2\sim x}-1)}{2\sim}\right\}\right] \quad (13)$$

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

For small interaction length, we may consider $2\sim x<1$. Consequently, we obtain

$$\left|\ (x)\right|^2=\left[\frac{|E_s(0)|}{|E_0(L)|}\right]^2\exp\left[2\{\sim(L-2x)+\copyright x\}\right] \qquad (14)$$

where $|E_s(0)|^2$ being the noise intensity for the SBS process [14] and its magnitude is generally taken to be about 10^{-12} to 10^{-13} times $|E_0(L)|^2$. Hence, significant TRR can be achieved in the crystal only if we find $2[\sim(L-2x)+\copyright x]\sim 30$ in equation (14) enabling oneself to have a gain $\sim e^{30}$ and reflectivity $\left|\ (x)\right|^2\sim 1$.

Since the time reversal reflection Brillouin mode is backscattered, it is easy to establish from Eq. (14) that at the entrance window ($x=L$) of the Brillouin cell of thickness L, we require $(\copyright-\sim)L\sim 15$ to achieve $\left|\ (x=L)\right|^2\sim 1$. This result can be compared very well with the observations of Zel'dovich et.al. [15] under the assumption of very low-loss Brillouin active media. From Eqs. (11) and (12), we find that SBS with finite gain could occur even at excitation intensities as low as to satisfy the threshold condition for Brillouin gain given by $\copyright\geq\sim$. But the above discussion makes it clear that excitation intensity $I_0(L)$ has to be much larger than the threshold value $I_{0,th}$ if one aims at the attainment of significant TRR-SBS reflectivity [i.e., $\left|\ (L)\right|^2\sim 1$ in a bulk crystal with $L=5$ mm.

It is clear from Eq. (12) that $I_{0,th}$ can be brought down by considering a semiconductor-laser interaction system yielding large i_B. Furthermore, the TRR-SBS reflectivity $\left|\ (L)\right|^2$ being directly dependent upon i_B, an enhancement in $\left|\ (L)\right|^2$ is also possible even at smaller sample length and not-too-high excitation intensity if we can achieve a large i_B in the crystal.

BRILLOUIN SUSCEPTIBILITY

We consider the well-known hydrodynamic model for a one-component (electron) semiconductor-plasma subjected to a pump electric field under thermal equilibrium. We have employed the coupled-mode scheme to obtain the nonlinear polarization with its origin being in the finite electrostrictive strain.

The basic equations employed in the formulation of i_B are as follows:

$$\frac{\partial^2 u}{\partial t^2}-\frac{C_a}{\circ}\frac{\partial^2 u}{\partial x^2}+2\Gamma_a\frac{\partial u}{\partial t}=\frac{¥}{2\circ}\frac{\partial}{\partial x}\left(E_e.E_1^*\right), \qquad (15)$$

$$\frac{\partial\vec{v}_0}{\partial t}+\neg\vec{v}_0=-\frac{e}{m}\left[\vec{E}_0+\left(\vec{v}_0\times\vec{B}_s\right)\right]=-\frac{e}{m}\vec{E}_e, \qquad (16)$$

$$\frac{\partial\vec{v}_1}{\partial t}+\left(\vec{v}_0\frac{\partial}{\partial x}\right)v_1+\neg\vec{v}_1=-\frac{e}{m}\left[\vec{E}_1+\left(\vec{v}_1\times\vec{B}_0\right)\right], \qquad (17)$$

$$\frac{\partial n_1}{\partial t}+v_0\frac{\partial n_1}{\partial x}+n_0\frac{\partial v_1}{\partial x}+D\frac{\partial^2 n_1}{\partial x^2}=0, \qquad (18)$$

$$P_{es}=-¥E_e\frac{\partial u^*}{\partial x}, \qquad (19)$$

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

$$\frac{\partial \vec{E}_x}{\partial x} = -\frac{n_1 e}{\pounds} + \frac{¥}{\pounds_1} E_0 \frac{\partial^2 u^*}{\partial x^2}. \qquad (20)$$

Eq. (15) represents the lattice motion in the crystal, where ° is its mass density, u is the lattice displacement, ¥ the electrostrictive coefficient, Γ_a the phenomenological damping parameter of acoustic mode and C_a the elastic constant. Eqs. (16) and (17) represent the zeroth and first-order oscillatory fluid velocities of an electron with effective mass '*m*' and charge '*e*' in which ¬ is the collision frequency. $\vec{E}_e$ represents the effective electric field which includes the Lorentz force $(\vec{v}_0 \times \vec{B}_0)$ in the presence of an external magnetic field $\vec{B}_0$. Eq. (18) 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. Eq. (19) reveals that the acoustic wave (AW) generated due to electrostrictive strain modulates the dielectric constant and gives rise to a nonlinear induced polarization P_{es}. The electric displacement in the presence of an external magnetic field is simply given by $\vec{D} = \pounds \vec{E}_{eff}$ [16]. The space charge field E_x is determined by the Poission's equation (20), where $\pounds_1$ is the dielectric constant of the crystal.

The interaction of the pump with the electrostrictively generated AW produces an electron density perturbation, which in turn derives 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 Singh and Aghamkar [17]. Differentiating Eq. (18) and using Eqs. (15) and (20), we obtain

$$\frac{\partial^2 n_1}{\partial t^2} + ¬ \frac{\partial n_1}{\partial t} + ¬D \frac{\partial^2 n_1}{\partial x^2} + ´ {}_p^2 n_1 + \frac{e k_s^2 n_0 ¥ u^* E_{eff}}{m \pounds_1} = i k_s n_1 \bar{E} \qquad (21)$$

where $\bar{E} = \frac{e}{m}\vec{E}_e$, $´ {}_p^2 = \left[µ_p^2 \left(\frac{¬^2}{¬^2 + µ_c^2} \right) \right]$, and $\vec{E}_e = \vec{E}_0 + (\vec{v}_0 \times \vec{B}_0)$.

Here $µ_c [= eB_s / m]$ is the electron cyclotron frequency and $µ_p [= (n_0 e^2 / m\pounds)^{1/2}]$ is the plasma frequency of carriers in the medium. We neglect the Doppler shift under the assumption that $µ_0 >> ¬ >> k_0 v_0$.

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_{fs}) corresponds to the first-order Stokes component of scattered light and varies as $\exp[i(k_s x - µ_s t)]$, whereas the slow component (n_{sl}) is associated with the AW and varies as $\exp[i(k_a x - µ_a t)]$.

We obtain coupled equations from Eq. (21) under rotating-wave-approximation [18]:

$$\frac{\partial^2 n_{fs}}{\partial t^2} + ¬ \frac{\partial n_{fs}}{\partial t} + ¬D \frac{\partial^2 n_{fs}}{\partial x^2} + ´ {}_p^2 n_{fs} + \frac{e k_s^2 n_0 ¥ u^* E_{eff}}{m \pounds_1} = -i k_s \bar{E} n_{sl}^* \qquad (22a)$$

and

$$\frac{\partial^2 n_{sl}}{\partial t^2} + ¬ \frac{\partial n_{sl}}{\partial t} + ¬D \frac{\partial^2 n_{sl}}{\partial x^2} + ´ {}_p^2 n_{sl} = -i k_s \bar{E} n_{fs}^* \qquad (22b)$$

From the above equations, it may be inferred that the generated AW and the Stokes mode couple to each other via the pump electric field in an electrostrictive medium. The slow component n_{sl} may be obtained from Eqs. (15) and (22) as:

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

$$n_{sl} = \frac{\varepsilon_0 k_s k_a n_0 \gamma^2 E_e E_1^* [A]^{-1}}{2\omega\varepsilon\left(\Omega_a^2 - 2i\Gamma_a\mu_a\right)} \tag{23}$$

where $A = [1-(\Omega_1^2 - i\mu_s\nu)(\Omega_2^2 + i\mu_a\nu)/k_s^2\overline{E}^2$,

$\Omega_a^2 = \mu_a^2 - k_a^2 v_a^2$, $\Omega_1^2 = \omega_p^2 - \mu_s^2 - k^2\nu D$, $\Omega_2^2 = \omega_p^2 - \mu_a^2 - k^2\nu D$.

The Stokes component of the nonlinear current density is obtained from the standard relation

$$J_{cd}(\mu_s) = n_{sl}^* e v_{0x} \tag{24}$$

The preceding analysis under RWA yields

$$J_{cd}(\mu_s) = -\frac{\mu_p^2\nu\varepsilon E_1}{(\nu^2 + \mu_c^2)} - \frac{\mu_p^2\varepsilon_0\gamma k_s k_a E_e^* E_0 E_s(\nu - \mu_0)[A]^{-1}}{2\omega\left(\Omega_a^2 - 2i\Gamma_a\mu_a\right)\left(\mu_c^2 - \mu_0^2\right)} \tag{25}$$

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

The induced polarization may be expressed as the time integral of the induced current density. The polarization $P_{cd}(\mu_s)$ may therefore be obtained from Eq. (25) as:

$$P_{cd}(\mu_s) = \frac{\mu_0^3 k_s k_a \mu_p^2 \varepsilon_0 |E_0|^2 E_s}{2\mu_s\omega\left(\mu_0^2 - \mu_c^2\right)^2\left(\Omega_a^2 - 2i\Gamma_a\mu_a\right)}[A]^{-1} \tag{26}$$

The origin of the SBS process lies in that component of $P_{cd}(\mu_s)$ which depends on $|E_0|^2 E_s$. Using the standard relation between induced polarization $P_{cd}(\mu_1)$ at frequency μ_s and Brillouin susceptibility $(\chi_B)_{cd}$, we obtain

$$(\chi_B)_{cd} = \frac{\mu_0^3 k_s k_a \mu_p^2 \gamma^2}{2\mu_s\omega\left(\mu_0^2 - \mu_c^2\right)^2\left(\Omega_a^2 - 2i\Gamma_a\mu_a\right)}[A]^{-1} \tag{27}$$

Besides $(\chi_B)_{cd}$, the system should also possess an electrostrictive polarization, which arises due to the interaction of the pump with the AW generated in the medium. The electrostrictive polarization is obtained from Eq. (19) as:

$$P_{es} = \frac{k_s k_a \gamma^2 \mu_0^4 |E_0|^2 E_s}{2\omega\left(\mu_0^2 - \mu_c^2\right)^2\left(\Omega_a^2 - 2i\Gamma_a\mu_a\right)} \tag{28}$$

The Brillouin susceptibility due to electrostrictive polarization is obtained as:

$$(\chi_B)_{es} = \frac{\varepsilon_0 k_s k_a \gamma^2 \mu_0^4}{2\omega\left(\mu_0^2 - \mu_c^2\right)^2\left(\Omega_a^2 - 2i\Gamma_a\mu_a\right)} \tag{29}$$

From Eqs. (27) and (29) we obtain the effective Brillouin susceptibility using the relation

$$\chi_B = (\chi_B)_{cd} + (\chi_B)_{es} = \frac{\varepsilon_0 k_s k_a \gamma^2 \mu_0^4\left(\Omega_a^2 + 2i\Gamma_a\mu_a\right)}{2\omega\left(\mu_0^2 - \mu_c^2\right)^2\left(\Omega_a^4 + 4\Gamma_a^2\mu_a^2\right)}\left[1 + \left(\frac{\mu_p^2}{\mu_0\mu_s}\right)[A]^{-1}\right] \tag{30}$$

Analyzing equation (30), we may find that the magnetic field (through μ_c) and high carrier density in diffusion driven semiconductors (through μ_p) affects χ_B.

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

RESULTS AND DISCUSSIONS

We now analyze numerically the nature of dependence of $\left| (x=L)\right|^2$ on different parameters such as excitation intensity, doping concentration (through ω_p), applied dc magnetic field (through ω_c) etc. The semiconductor crystal used for this purpose is n-InSb at 77 K duly irradiated by a nanosecond pulsed 10.6 μm CO_2 laser. The physical constants used are [17]: $m = 0.014 m_0$ (m_0; free-electron mass), $\varepsilon_l = 15.8$, $\mu = 5\times10^{-10}\, Fm^{-1}$, $\rho = 5.8\times10^{9}\, Kgm^{-3}$, $\tau = 3.5\times10^{11}\,\text{sec}^{-1}$, $\Gamma_a = 5\times10^{8}\,\text{sec}^{-1}$, $\mu_a = 4.8\times10^{9}\,\text{sec}^{-1}$, $k = 6\times10^{5}\, m^{-1}$, $\sigma = 50 m^{-1}$, and $\mu = 1.78\times10^{14}\,\text{sec}^{-1}$.

The threshold value of pump intensity required for the onset of time reversal reflection via SBS process is obtained in magnetized semiconductors and is given as:

$$I_{0,th} = \frac{\left|\varepsilon_0 c^3 k\sigma\left[2\rho\left(\mu_0^2-\mu_c^2\right)^2\left(\omega_a^4+4\Gamma_a^2\mu_a^2\right)\right]\right.}{\mu^2\varepsilon_0 k_s k_a \mu^2 \mu_0^4\left(\omega_a^2+2i\Gamma_a\mu_a\right)\left[1+\left(\mu_p^2/\mu_0\mu_s\right)[A]^{-1}\right]} \qquad (31)$$

At pump intensities $I_0(L) > I_{0,th}$, one can get finite $\left|(L)\right|^2$. Eq. (31) manifests the critical dependence of $I_{0,th}$ on the applied dc magnetic field (via μ_c) and the carrier concentration n_0 (via μ_p). The characteristic dependence of threshold intensity $I_{0,th}$ on doping concentration (in terms of μ_p/μ_0) with $\mu_c = 0$ and $\mu_c = \mu_0$ is shown in Fig. 1.

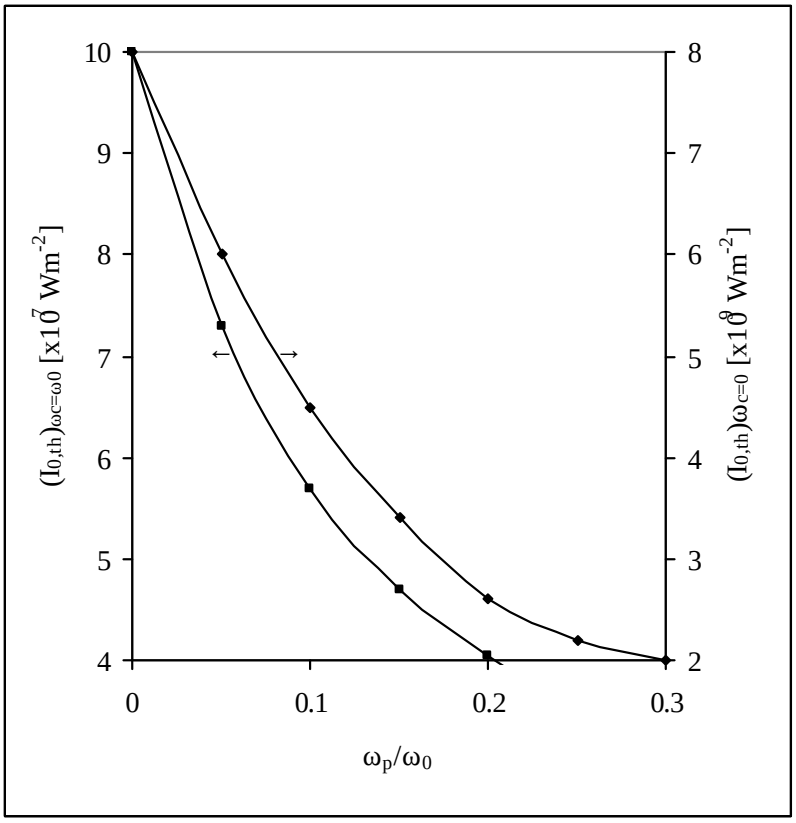

Fig. 1. Variation of threshold pump intensity I_{0th} of the TRR-SBS process with doping concentration (in terms of μ_p/μ_0) in n-InSb crystal with $\mu_c = \mu_0$ and $\mu_c = 0$.

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

In both the cases $I_{0,th}$ decreases with increase in doping concentration and saturates at higher values of doping concentration. In absence of magnetic field ($\mu_c = 0$), $I_{0,th}$ starts with $8\times10^{9} Wm^{-2}$ at $\mu_p = \mu_0$ and saturates at lower value $2\times10^{9}\ \mathrm{Wm}^{-2}$ at $\mu_p/\mu_0 = 0.3$. In presence of magnetic field (when $\mu_c = \mu_0$), $I_{0,th}$ starts with $1\times10^{8} Wm^{-2}$ at $\mu_p = \mu_0$ and saturates at lower value $4\times10^{7} Wm^{-2}$ at $\mu_p/\mu_0 = 0.3$. A comparison between both cases shows that an externally applied magnetic field considerably reduces the threshold field for TRR-SBS.

Fig. 2 shows the variation of $I_{0,th}$ of the TRR-SBS process with magnetic field (in terms of μ_c/μ_0) in n-InSb crystal. It can be seen that $I_{0,th}$ is fairly independent on magnetic field when $\mu_c < \mu_0$. Further increase in the value of magnetic field causes sharp fall in $I_{0,th}$. Around $\mu_c \sim \mu_0$, $I_{0,th}$ attains a minimum value $\sim 5.2\times10^{7} Wm^{-2}$ when the crystal is irradiated by CO_2 laser of frequency $1.78\times10^{14} s^{-1}$.

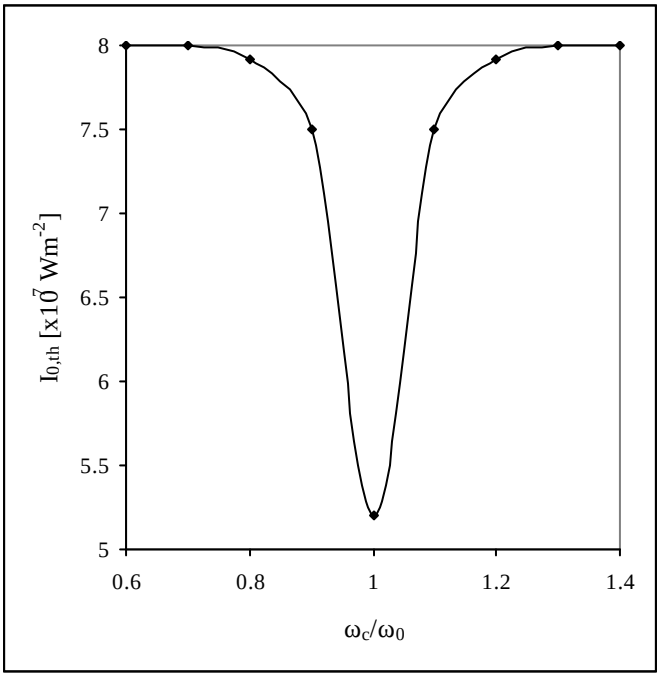

Fig. 2. Variation of threshold pump intensity I_{0th} of the TRR-SBS process with magnetic field (in terms of μ_c/μ_0) in n-InSb crystal

For pump intensity $I_0(L) > I_{0,th}$, one can obtain finite $|\ (L)|^2$ at the entrance window by using equation (14) as:

$$|\ (L)|^2 = \left[\frac{|E_s(0)|}{|E_0(L)|}\right]^2 \exp[2(\copyright - \sim)L] \qquad (32)$$

At $I_0(L) > I_{0,th}$, $\copyright > \sim$, but since the ratio $|E_s(0)/E_0(L)|^2$ is usually taken to be as small as 10^{-13}, considerably large TRR-SBS reflectivity (i.e., $|\ (L)|^2 \sim 1$) can be achieved when

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

$(© - \sim)L \cong 15$ (33)

such that $\exp[2(© - \sim)L] \sim 10^{13}$. We may find such large reflectivity without much difficulty in a material system such that is in the form of low-loss optical fibre or waveguide having very large interaction path length and which is duly irradiated by a not-too-intense off-resonant laser. As is well known in the presence of a large magnetic field, a III-V semiconductor like InSb can exhibit giant $¡_B$, an interaction length in the mm range can be sufficient to get $\left| (L)\right|^2 \sim 1$ at pump intensities below the optical damage threshold. Accordingly, we consider a bulk InSb crystal with $L = 5$ mm subjected to off-resonant nanosecond pulsed 10.6 μm CO_2 laser excitation.

The characteristic dependence of $\left| (L)\right|^2$ on the pump intensity, excess carrier concentration, and the magnetic field in n- InSb are plotted in Figs. 3-5.

Fig. 3 illustrates the nature of dependence of $\left| (L)\right|^2$ on the pump intensity $I_0(L)$ in the presence ($\mu_c = \mu_0$) and absence ($\mu_c = 0$) of applied magnetic field. In both the cases, $\left| (L)\right|^2$ increases with increase in pump intensity. Comparison of both cases reveals that $\left| (L)\right|^2$ in presence of magnetic field (when μ_c is tuned with μ_0) is about hundred times larger than in the absence of magnetic field.

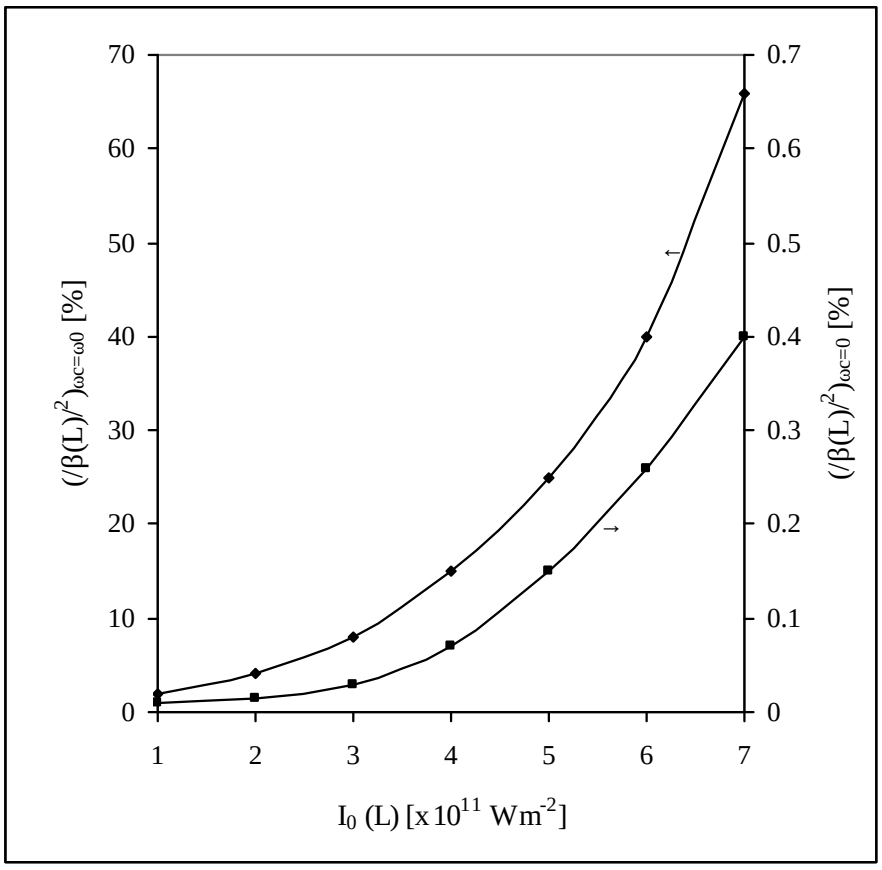

Fig. 3. Dependence of TRR-SBS reflectivity $|\beta(L)|^2$ on the pump intensity $I_0(L)$ in n-type InSb crystal at 77 K with $\mu_c = \mu_0$ and $\mu_c = 0$.

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

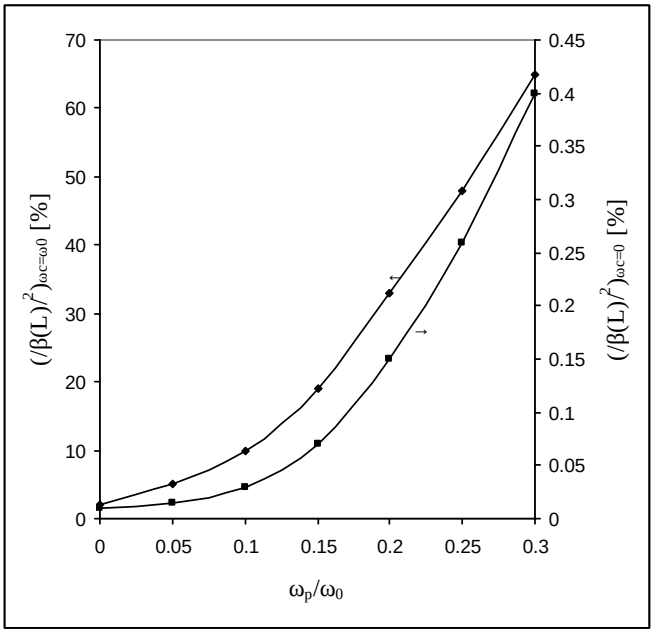

Fig. 4. Variation of TRR-SBS reflectivity $|\beta(L)|^2$ with the electron concentration (in terms of μ_p/μ_0 in InSb crystal with $\mu_c = \mu_0$ and $\mu_c = 0$ at $I_0 = 7\times10^{11}\,\mathrm{Wm}^{-2}$.

Fig. 4 shows the variation of $|(L)|^2$ with the electron concentration (in terms of μ_p/μ_0) in InSb crystal with $\mu_c = 0$ and $\mu_c \sim \mu_0$, at $I_0 = 7\times10^{11} Wm^{-2}$. In both the cases $|(L)|^2$ increases rapidly as the carrier density increases. The presence of magnetic field (when $\mu_c \sim \mu_0$) enhances the reflectivity by a factor of 10^2.

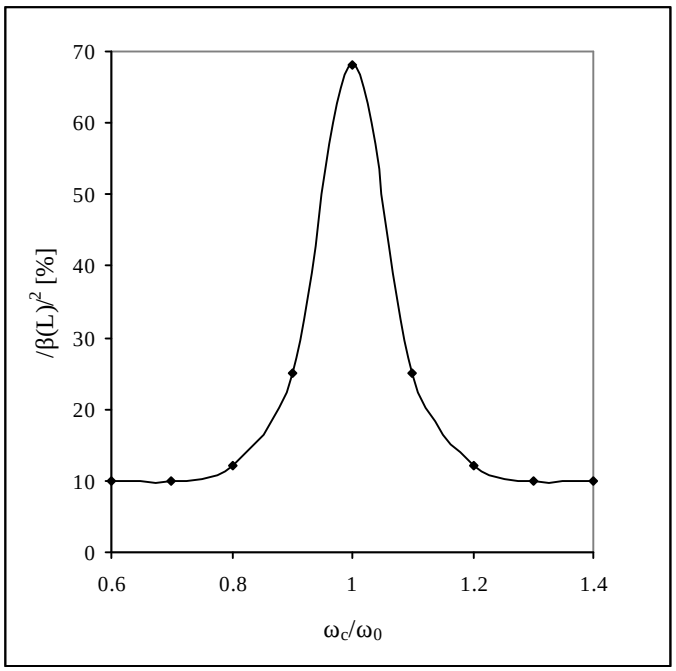

Fig. 5. Dependence of TRR-SBS reflectivity $|\beta(L)|^2$ on magnetic field (in terms of μ_c/μ_0) in InSb crystal with $I_0 = 7\times10^{11}\,\mathrm{Wm}^{-2}$.

International Journal of Engineering Technology Science and Research
IJETSR
www.ijetsr.com
ISSN 2394 – 3386
Volume 4, Issue 7
July 2017

Fig. 5 shows the dependence of $|(L)|^2$ on magnetic field (in terms of μ_c/μ_0) in InSb crystal with $I_0 = 7\times10^{11} Wm^{-2}$. It is a unique feature displayed in this figure that finite reflectivity is obtained only when μ_c is very close to μ_0. This behaviour may be utilized for the construction of magnetic switches. One can infer from this figure that reflectivity increases with magnetic field up to $\mu_c \le \mu_0$. The reflectivity starts decreasing in the region when $\mu_c \ge \mu_0$. This behaviour may be attributed to the fact that i_B is proportional to $\left(\mu_0^2 - \mu_c^2\right)^{-2}$ (Eq. 31), and when $\mu_c > \mu_0$ magnetic absorption becomes vital which reduces the TRR reflectivity. Thus $|(L)|^2 \approx 70\%$ is obtained in magnetized semiconductors when electron-cyclotron frequency is tuned with pump frequency.

From the above discussion, we conclude that the magnetized diffusion driven semiconductors may be the active materials for fabrication of SBS based time reversal reflection mirrors with high reflectivity.

REFERENCES

[1] Yariv, A., Yeh, P., 1984 “Optical Waves in Crystals”, Wiley, New York, pp. 549-62.
[2] Grower, M., Proch, D., 1994 “Optical Phase Conjugation”, Springer, Berlin.
[3] Brignon, A., Huignard, J.P., 2003 “Phase Conjugate Laser Optics” Wiley, New York.
[4] Jackel, S., Moshe, I., Lavi, R., 2003 “Comparison of adaptive optics and phase conjugate mirrors for correction of aberrations in double-pass amplifiers”, Appl. Opt. 42, 983-989.
[5] Kong, H.J., Yoon, J.W., Shin, J.S., Beak, D.H., Lee, B.J., 2006 “Long term stabilization of the beam combination laser with a phase controlled stimulated Brillouin scattering phase conjugation mirrors for the laser fusion driver”, Laser Part. Beams 24, 519-523.
[6] Pepper, D.M., 1985 “Laser Handbook”, ed. Stitch, M.L., Bass, M., Elsevier, Amsterdam, pp. 333-467.
[7] Shakir, S.A., 1985 “Increasing the efficiency of stimulated scattering phase conjugate mirrors” Proc. SPIE 540, Southwest Conf. on Optics ’85, 303.
[8] Singh, M., Aghamkar, P., Duhan, S., 2008 “Enhancement of second- and third-order optical susceptibilities in magnetized semiconductors”, Chin. Phys. Lett. 25, 3276-3279.
[9] Pepper, D.M., 1982 “Nonlinear optical phase conjugation” Opt. Engg. 21, 155-158.
[10] Sen, P.K., Sen, P., 1988 “Optical phase conjugation in narrow direct gap semiconductors”, J. Appl. Phys. 63, 4637-4643.
[11] Yeh, P., 1989 “2-wave mixing in nonlinear media”, IEEE J. Quantum Electron. QE-25, 484-519.
[12] Pepper, D.M., 1985 in Laser Handbook, ed. Stitch, M.L., Bass, M., Elsevier, Amsterdam, vol. 4, pp. 333-485.
[13] Dinu, M., 2003 “Dispersion of phonon-assisted non-resonant third-order nonlinearities”, IEEE J. Quantum Electron. QE-39, 1498-1503.
[14] Zel’dovich, B. Ya., Popovichev, V.I., Ragul’skii, V.V., Faizullov, F.S. 1972 “Connection between the wavefronts of the reflected and exciting light in stimulated Mandel’shtam – Brillouin scattering”, JETP Lett. 15, 109-113.
[15] Zeldovich, B. Ya., Pilipetsky, N.F., Shkunov, V.V., 1983 “Optical Phase Conjugation”, ed. Fisher, R.A., Academic, New York, pp. 135-167.
[16] Landau, L.D., Lifshitz, E.M., 1963 “Electrodynamics of Continuous Media”, Peragmon Press, Oxfors, pp. 337-340.
[17] Singh M., Aghamkar, P., 2008 “Coherent Brillouin scattering in noncentrosymmetric semiconductors”, J. Mod. Opt. 55, 1251-1265.
[18] Von der Lind, D., Maier, M., Kaiser, W., 1969 “Quantitative investigations of the stimulated Raman effect using sub-nanosecond light pulses”, Phys. Rev. 178, 11-17.