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Stimulated Raman Scattering: Temperature and Band Structure Effects

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Abstract

Including the hot carriers and band structure modifications, in this paper authors present an analytical investigation of steady-state and transient gain characteristics of the Raman scattered Stoke mode resulting from the nonlinear interaction of an intense pumping light beam with molecular vibrations of semiconductor plasma. The origin of this nonlinear interaction lies in the third order (Raman) susceptibility arising from the induced current density and density fluctuations generated within the medium. The threshold intensities, steady state gain and transient gain are determined through the effective susceptibility derived with the coupled mode theory of plasmas under hydrodynamic regime. The numerical estimates are made for CdS, GaAs and InSb crystals duly irradiated by 10.6 μm CO_2 laser. The analysis establishes that the crystal with higher density of state mass proves its potential as candidate material for the fabrication of cubic nonlinear devices.

Key Words: Direct Band Gap Semiconductors, Stimulated Raman Scattering, Nonlinear Amplification, Optical Transient Phenomenon.

1. Introduction

In a nonlinear optical mixing process, at large amplitude of the incident electromagnetic waves, a number of nonlinear mode-coupling interactions may take place which modify the electrical and optical properties of the medium. The optical properties of a medium can also be modified by an externally applied electric field [1]. Out of the several nonlinear mode coupling interactions, Stimulated Raman scattering and consequent Raman gain occupies a distinct place. There are numerous report on the Stimulated Raman scattering in semiconductor medium during the last two decades viz. Dubey and Ghosh [2] and Singh et al. [3]. They have reported the nonlinear absorption and refractive index of a Raman scattered mode in

magnetoactive centrosymmetric semiconductor plasmas and pointed out that significant Raman growth rate is possible if the pump field strength is very large ($\sim 10^8 Vm^{-1}; corresponding\ \text{int}ensity \sim 5.3\times 10^{13} Wm^{-2}$) but silent for regarding the pulse duration of the pump wave. It is established fact that at such a high intensity pump irradiation the free carriers of the medium gain energy through collisions with the lattice. During this process of gaining energy, the temperature of the free carriers do not remain at lattice temperature but it increases with pump intensity and subsequently alters several physical parameters responsible for Stimulated Raman Scattering to occur.

Forchel et al. [4] have derived a very useful algebric expression for maximum possible plasma density (free carrier concentration) as a function of carrier temperature and density of state mass (function of band structure) of the crystal.

From the available literature and as per the knowledge of the present authors, nobody has studied the effect of carrier temperature and density of state mass on the complete gain characteristics of stimulated Raman scattering in direct gap semiconductors.

Motivated by the present state of art the authors have investigated the stimulated Raman scattering incorporating band structure and hot carrier effects in direct gap semiconductors. The inclusion of density of states mass through band structure and carrier temperature adds new dimensions to the present study. Initially the crystal temperature is assumed to be maintained at room temperature (300K), therefore the transfer of energy and momentum of carriers occurs through their collisions with polar optical phonons (POP) and acoustic phonons (AP) respectively. In the present case, the nonlinearity is solely due to the induced nonlinear current density of free carriers. However, as far as we know, no such attempt has been made to determine steady-state and transient gain coefficients of scattered wave in direct gap semiconductor plasmas incorporating the density of states mass of free carriers.

2. Theoretical Formulation

In order to study the stimulated Raman scattering in direct band gap semiconductor plasmas arising due to third-order susceptibility $\chi^{(3)}$, the hydrodynamic model of the homogeneous one component (n-type) semiconductor plasmas is considered, satisfying the condition $kl << 1$. This condition implies that the sound wave length ($\lambda = 2\pi / k$) is much greater than the mean free path (l) of electrons, so that the motion of carriers under the influence of the external fields is averaged out. We have used the particular geometry where the incident high frequency spatially uniform laser radiation (pump wave) $E_0 \exp(-i\omega_0 t)$ is applied parallel to the propagation vector k (along the x-axis). The three direct band gap semiconductor crystals (InSb, GaAs and CdS) are assumed to possess an isotropic and nondegenerate conduction band and being centrosymmetric in nature the effect of any pseudopotential is neglected for simplicity. In present analysis the Raman medium is taken as consisting of N harmonic molecular oscillators per unit volume; each oscillator being characterized by

its position x and the normal vibrational coordinate $u(x,t)$.

The optical-phonon mode is presented in an one dimensional configuration as [5].

$$\frac{\partial^2 u(x,t)}{\partial t^2} + \Gamma \frac{\partial u(x,t)}{\partial t} + \omega^2{}_T\, u(x,t) = \frac{F(x,t)}{M} \tag{1}$$

where Γ is the damping constant corresponding to the phenomenological phonon collision frequency $(\approx 10^{-2}\,\omega_T)$ [5], ω_T being the undamped molecular vibrational frequency and is considered to be equal to the transverse optical phonon frequency. $F(x,t)$ is the driving force per unit volume, which can be obtained by considering the electromagnetic energy in the presence of the molecules and in polarisable material is given by

$$F(x,t) = \frac{1}{2}\varepsilon(\partial\alpha/\partial u)_0 \overline{E}^2(x,t), \tag{2}$$

where $\varepsilon = \varepsilon_0\varepsilon_\infty$, ε_0 and ε_∞ being the absolute permittivity and the high frequency permittivity respectively; $(\partial\alpha/\partial u)_0$ is the differential polarizability and the bar over any quantity indicates the averaging over a few optical periods. This shows that the molecular vibrations within the medium can be driven by the electric field due to nonvanishing differential polarizability $(\partial\alpha/\partial u)_0$. The other fundamental equations employed are given in one dimension as

$$\frac{\partial \vartheta_0}{\partial t} + \upsilon \vartheta_0 = \frac{e}{m} E_0, \tag{3}$$

$$\frac{\partial \vartheta_1}{\partial t} + \upsilon \vartheta_1 + (\vartheta_0 \frac{\partial}{\partial x})\vartheta_1 = \frac{e}{m} E_1 \tag{4}$$

$$\frac{\partial n_1}{\partial t} + \vartheta_0 \frac{\partial n_1}{\partial x} + (n_0 \frac{\partial \vartheta_1}{\partial x}) = 0 \tag{5}$$

$$P = \varepsilon N(\partial\alpha/\partial u)_0 u^* E_0 \tag{6}$$

$$\frac{\partial E_1}{\partial x} = \frac{n_1 e}{\varepsilon} - N(\partial\alpha/\partial u)_0 E_0 (\partial u^*/\partial x) \tag{7}$$

Equation (3) is the momentum transfer equation in which ϑ_0 is the equilibrium electron velocity and shows that the electrons will oscillate under the influence of the applied time varying pump electric field having effective mass m and charge e. υ is the phenomenological electron collision frequency. Similarly eq. (4) represents the first order equation of motion for the scattered wave. Equation (5) represents the equation of continuity in which n_0 and n_1 are the equilibrium and perturbed carrier densities, respectively. Equation (6) represents the nonlinear polarization of the medium. The space charge field E_1 is determined by the Poisson equation (7) in which the second term represents the polarization due to the molecular vibrations driven by the electric field. These one dimensional equations are appropriate for the nondegenerate semiconductors. The molecular vibrations at frequency ω modulates the effective dielectric constant of the medium leading to the exchange of energy

among the electromagnetic fields separated in frequency by integral multiples of ω [i. e. $(\omega_0 \pm p\omega)$ where $p = 1, 2, ...$]. The polarization of the medium at very high frequencies is considered on neglecting the interaction of the electrons with one another and the nuclei of the atoms as the impressed frequencies of the field are fairly large as compared to the frequencies of the motion of the electrons in the medium. Thus the electric induction is given by $\vec{D} = \varepsilon \vec{E} + \vec{P}$ [6]. In the above expressions the effect of $(\vec{\vartheta_0} \times \vec{B_1})$ term is neglected by assuming that the wave propagates along such a direction of the crystal which produces a longitudinal electric field.

2.1. Heating of Carriers and Temperature Dependent Plasma Density

Now to compute the carrier temperature, using Eq. (3) and (4) the perturbed electron fluid velocities due to pump and scattered waves may be deduced as

$$\vartheta_0 = \frac{eE_0}{m(\nu - i\omega_0)} \tag{8a}$$

and

$$\vartheta_1 = \frac{eE_1}{m(\nu - i\omega_1)} \tag{8b}$$

It is worth mentioning here that the fundamental requirement for the SRS to occur is that the applied electric field should be well above defined pump amplitude known as threshold pump amplitude. When this high intensity pump field interacts with a high mobility semiconductor, carriers acquire momentum and energy from the incident wave and consequently acquire a temperature (T_e) somewhat higher than that of lattice temperature (T_0).

Following Sodha et al. [7] and using eq. (8a) for the said geometry, the time independent part of power absorbed per electron from pump is

$$\frac{e}{2}\mathrm{Re}(\nu_0 E_0^{\,*}) = \frac{\nu e^2 E_0 E_0^{\,*}}{2m\left(\nu^2 + \omega_0^{\,2}\right)} \tag{9}$$

in which " * " and " Re " denote the complex conjugate and real part of the quantity, respectively

Following Conwell [8], the power dissipation per electron in collisions with the polar optical phonons (POP) may be expressed as

$$\langle p \rangle_{pop} = \left(\frac{2K_B\theta_D}{m\pi}\right)^{1/2} eE_{PO}x_e^{\,1/2}K_0\left(\frac{x_e}{2}\right)\exp\left(\frac{x_e}{2}\right)\frac{\exp(x_0 - x_e) - 1}{\exp(x_0) - 1} \tag{10}$$

where $x_{0,e} = \dfrac{\hbar\omega_1}{K_B T_{0,e}}$ in which $\hbar\omega_1$ is the energy of POP given by $\hbar\omega_1 = K_B\theta_D$. θ_D is

the Debye temperature of the medium. $E_{PO} = \frac{mc\hbar\omega_0}{\hbar^2}\left(\frac{1}{\varepsilon_\infty} - \frac{1}{\varepsilon_L}\right)$ is the field of POP scattering potential in which ε_L and ε_∞ are the lattice and high frequency dielectric permittivity of the medium, respectively, $K_0\left(\frac{x_e}{2}\right)$ is the zeroth-order Bessel function of first kind.

In steady state, the power absorbed per carrier from the pump electric field is just equal to the power dissipated per carrier due to collision with POP scattering. Therefore, for moderate heating of carriers by pump electric field, we obtain

$$\frac{T_e}{T_0} = 1 + \frac{e^2\nu}{2m}\frac{\tau E_0 E_0^*}{\left(\nu^2 + \omega_0^2\right)} \tag{11}$$

Now with an aim to include the temperature and band structure dependency of plasma density in our calculation we shall follow the work of Forchel et al. [4]. They have developed an empirical relation for maximum plasma density n_0 in direct gap semiconductor as a function of band structure (density of states mass) and carrier temperature, which is given as

$$n = 4.2\times10^{22}\left(\frac{m_{de}\times m_{dh}}{m_{de}+m_{dh}}\right)^{\frac{3}{2}} T^{3/2} \tag{12}$$

in which m_{de} and m_{dh} are the density of states mass of electrons and holes in the medium.

Physically this relation means that the density maximum is related to the constant degeneracy of the plasma. The relation holds good for effective degeneracy's close to two, for which the plasma transport is significant, and is affected by both, temperature and density gradients. They have reported that experimentally observed band structure dependency of the reduced densities in direct gap semiconductors is same as the band structure variation of the non equilibrium electron hole plasma observed in the indirect gap materials. The induced plasma densities are in good agreement with their productions using the thermo-diffusion model.

The high frequency pump field gives rise to a carrier density perturbation within the Raman active medium. In a nondegenerate semiconductors these density perturbations can be obtained by using the standard approach [9].

Using eqs. (3) - (6), we obtain after some algebric simplification the following relation

$$\frac{\partial^2 n_1}{\partial t^2} + \upsilon\frac{\partial n_1}{\partial t} + \omega_R^2 n_1 - n_0 N(\partial\alpha/\partial u)_0 E_0(\partial u^*/\partial x) = -\overline{E}\frac{\partial n_1}{\partial x}, \tag{13}$$

where,

$$\overline{E} = \frac{e}{m}E_0;\ \omega_R{}^2 = \frac{\omega_p{}^2\omega_L{}^2}{\omega_T{}^2};\ \frac{\omega_L}{\omega_T} = \sqrt{\frac{\varepsilon_L}{\varepsilon_\infty}};\ \omega_p{}^2 = \frac{n_0 e^2}{m\varepsilon_0\varepsilon_L}$$

Here ω_L is the longitudinal optical phonon frequency and is given by $\omega_L = k_B\theta_D / h$, where k_B and θ_D are the Boltzman constant and Debye temperature of the crystal, respectively.

In deducing eq. (13) we have neglected the Doppler shift under the assumption $\omega_0 \rangle\rangle \upsilon \rangle\rangle k\vartheta_0$. The spatial dependency of the pump has also been neglected as the laser irradiation is assumed to be uniform in space compared to the product modes [i. e. $|k_0| \langle\langle |k|$]. The scattered Raman mode oscillates at forced electromagnetic frequencies $(\omega_0 \pm \omega)$; the higher order terms with frequencies $(\omega_0 \pm p\omega)$ [p=2, 3….] being off resonant are neglected. In the present paper we have considered only the Stokes component of the scattered electromagnetic wave $\omega_s = \omega_0 - \omega$ is obtained as,

$$u = \frac{(\varepsilon / 2M)(\partial\alpha / \partial u)_0 E_0 E_1{}^*}{(\omega_T{}^2 - \omega^2 - i\Gamma\omega)} \tag{14}$$

where $E_1{}^*$ represents the complex conjugate of the scattered polarized electromagnetic wave amplitude.

The perturbed electron density (n_T) of the Raman active medium due to molecular vibrations can be deduced using eqs. (7) and (14) as,

$$n_T = \frac{ik \in}{e}\left[\frac{\left(\omega_T{}^2 - \omega^2 + i\Gamma\omega\right) + \left(\frac{\in N}{2M}\right)\left(\frac{\partial\alpha}{\partial u}\right)_0{}^2 |E_0|^2}{\left(\frac{\in N}{2M}\right)\left(\frac{\partial\alpha}{\partial u}\right)_0 E_0{}^*}\right] u^* \tag{15}$$

The density perturbations associated with the molecular vibrations at frequency (ω) beats with the pump at frequency (ω_0) and produce side band frequencies. We consider only the Stokes mode at frequency $(\omega_0 - \omega)$. The slow component of the density perturbation may be obtained from eq (13) as

$$n_s = \frac{-ikeE_0 n_T{}^*}{m\left(\delta_1{}^2 - i\omega_s\nu\right)} \tag{16}$$

Here $\delta_1{}^2 = (\omega_R{}^2 - \omega_S{}^2)$ and suffixes T and s denotes the component of the perturbed electron density associated with the molecular vibrations and Stokes mode respectively.

where $\omega_R{}^2 = \dfrac{\omega_p{}^2\omega_L{}^2}{\omega_T{}^2}$

The propagation of scattered Stokes component of the Raman mode is governed

by the general wave equation obtained from eq. (6). Now $\vec{J}$ is the resonant Stokes component of the current density due to finite nonlinear induced polarization of the medium that can be represented as,

$$J(\omega_s) = n_0 e \vartheta_1 + n_s^{*} e \vartheta_0 \tag{17}$$

The preceding analysis yields,

$$J(\omega_s) = \frac{\omega_p^{\,2} \varepsilon E_0}{(\upsilon - i\omega_0)} - \frac{\varepsilon k^2 |E_0|^2 E_1}{(\delta_1^{\,2} + i\omega_s \upsilon)} \times [1 - \frac{(\varepsilon N / 2M)(\partial\alpha / \partial u)_0^{\,2} |E_0|^2}{(\delta_2^{\,2} + i\Gamma\omega)}] \tag{18}$$

where $\delta_2^{\,2} = (\omega_T^{\,2} - \omega_S^{\,2})$

The first part of eq. (17) represents the linear component of the induced current density while the latter term represents the nonlinear coupling amongst the three interacting waves via the nonlinear current density $J_{nl}(\omega_s)$. In this present investigation the effect of the transition dipole moment is neglected in order to study the effect of nonlinear current density on the induced polarization in a Raman active medium.

Henceforth, treating the induced polarization P_{cd} as the time integral of the current density J_{nl} we may write

$$P_{cd}(\omega_s) = \int J_{nl}(\omega_s)\, dt \tag{19}$$

We obtain the nonlinear induced polarization due to perturbed current density from eqs. (18) and (19) as

$$P_{cd}(\omega_s) = \frac{-\varepsilon_0 \varepsilon_\infty k^2 e^2 |E_0|^2 E_1}{\omega_s m^2 \omega_0 (\delta_1^{\,2} + i\omega_s \upsilon)} [1 - \frac{(\varepsilon N / 2M)(\partial\alpha / \partial u)_0^{\,2} |E_0|^2}{(\delta_2^{\,2} + i\Gamma\omega)}] \tag{20}$$

The induced polarization at the Stokes frequency ω_s is defined by

$$P_{cd}(\omega_s) = \varepsilon_0 \chi_R^{(3)} |E_0|^2 E_1(\omega_s) \tag{21}$$

It is well known that the origin of the SRS process lies in that component of $P_{cd}(\omega_s)$ which depends on $|E_0|^2 E_1$ and consequently depend on third order susceptibility of the Raman active medium which is known as the Raman susceptibility χ_R that can be obtained using eqs. (20) and (21) as

$$[\chi_R^{(3)}]_{cd} = \frac{-\varepsilon_\infty k^2 e^2}{\omega_s m^2 \omega_0 (\delta_1^{\,2} + i\omega_s \upsilon)} \tag{22}$$

Besides Raman susceptibility, the system should also possess a polarization induced by the interaction of the pump wave with the molecular vibrations generated in the medium such as,

$$P_{mv}(\omega_s) = \varepsilon_0 [\chi_R^{(3)}]_{mv} |E_0|^2 E_1 \tag{23}$$

Using eqs. (6) and (23), we obtain

$$[\chi_R^{(3)}]_{mv} = \frac{\varepsilon_\infty (\varepsilon N / 2M)(\partial\alpha / \partial u)_0^{\,2}}{(\delta_2^{\,2} + i\Gamma\omega)} \tag{24}$$

Thus the effective Raman susceptibility of the centrosymmetric medium is obtained as

$$[\chi_R^{(3)}]_{eff} = [\chi_R^{(3)}]_{mv} + [\chi_R^{(3)}]_{cd}$$
$$= \frac{\varepsilon_\infty (\varepsilon N / 2M)(\partial\alpha / \partial u)_0^{\,2}}{(\delta_2^{\,2} + i\Gamma\omega)} + \frac{-\varepsilon_\infty k^2 e^2}{\omega_s m^2 \omega_0 (\delta_1^{\,2} + i\omega_s \upsilon)} \tag{25}$$

As the noncentrosymmetricity of the medium induces a retardation effect on Raman susceptibility due to finite $\chi^{(2)}$ we have chosen direct band gap crystals possessing inversion symmetry to derive eq. (25). It can be observed from this equation that the effective susceptibility is strongly depend on the differential polarizability of the medium.

We now estimate the Raman gain constant $g_R(\omega_s)$ of the Stokes mode which is related to the imaginary part of $[\chi_R^{(3)}] = [\chi_{Rr}^{(3)}] + [\chi_{Ri}^{(3)}]$ through,

$$g_R(\omega_s) = \frac{-\omega_s}{\eta c} [\chi_{Ri}^{(3)}] |E_0|^2 \tag{26}$$

where η is the refractive index of the medium and c is the speed of light in the medium. In order to attain a growth of the signal $(g_R \rangle 0)$, it may be infer that $\left(\chi_{Ri}^{(3)}\right)$ should be negative.

The threshold value of the pump amplitude required for the onset of the SRS process is obtained from eq. (25) as

$$E_{th} = \frac{(\delta_2^{\,2} + i\Gamma\omega)^{1/2}}{(\varepsilon N / 2M)^{1/2} (\partial\alpha / \partial u)_0}, \tag{27}$$

It can be observed from eq. (27) that the Raman instability of the signal wave has a nonzero intensity threshold.

2.2 Transient gain coefficient

In addition to the steady state Raman gain coefficient, it is found that the transient gain coefficient is of great significance not only for predicting the threshold pump intensity for the onset of the instability exactly but also in predicting the optimum pulse durations for which the stimulated Raman scattering process can be observed. Therefore, we will extend the above formulations to study the transient behaviour of the nonlinear medium. To do so, we consider the pump pulse duration $\tau_p \leq$ acoustic

phonon life time (τ_a) and following Carman et al. [10], we obtain

$$g_{TR} = [2g_{TR}x\Gamma\tau_p]^{\frac{1}{2}} - \Gamma\tau_p \tag{28}$$

where Γ being the optical phonon life time, x is the interaction length. Here, it is worth mentioning that the interaction length (x) becomes very small for modulated side band mode because the side band pulse and the laser pulse travel in different directions and hence, their overlap region cannot exceed the length of the laser pulse; viz., for a typical picoseconds pulse laser, the interaction length is of the order of a millimeter. From all practical point of view, the interaction length should be replaced by $\left(c_1 c_p / 2\right)$ (where c_1 is the velocity of light in the crystal medium) when the pump and the signal wave durations are shorter than 10 ps [11]. Consequently the threshold pump intensity for the onset of transient Raman scattering can be obtained by making $g_{TR} = 0$ in Eq. (28). This yields

$$I_{th} = \frac{\Gamma}{2Gc_1} \tag{29}$$

with $G = \frac{g_R}{I_{in}}$, the steady-state gain per unit pump intensity.

Here one may notice that for a laser pulse as long as twice the acoustic-phonon lifetime, the threshold intensity for SRS process is raised by a factor of two.

For comparatively long pulse-duration $\left(\tau_p \geq 10^{-10} s\right)$, the cell length can be taken within the range $\left(10^{-4} m - 10^{-6} m\right)$, the cell length can be taken equal to (x) and under such circumstances, we find

$$g_{TR} = \left(\Gamma\tau_p\right)^{\frac{1}{2}}[-\left(\Gamma\tau_p\right)^{\frac{1}{2}} + \left(g_R x\right)^{\frac{1}{2}} \tag{30}$$

From the above expression one may obtain the cut-off value of the pulse duration $\left(\tau_p\right)_{opt}$, above which no transient gain could be achieved. This can be obtained by making $g_{TR} = 0$, which yields

$$\left(\tau_p\right)_{opt} \approx \left(\frac{g_R x}{\Gamma}\right) \tag{31}$$

Eq. (30) reveals that for a unit interaction length, the optimum pulse depends on the ratio of the steady state gain coefficient and the acoustic phonon life time.

Hence we may obtain the order of magnitude of optimum pulse duration for different media

For CdS $\left(\tau_p\right)_{opt} = (3.941\times10^{-5} I_{in})\, s$

GaAs $\left(\tau_p\right)_{opt} = (3.793\times10^{-6} I_{in})\, s$

InSb $\left(\tau_p\right)_{opt} = (2.823\times10^{-7} I_{in})\, s$

These values of $(\tau_p)_{opt}$ not only explains the washing out of gain of scattered wave at large pulse durations but also suggests that optimum pulse duration can be increased by increasing the pump intensity.

3. Results and Discussion

The analytical investigation for the possibility of stimulated Raman scattering and the consequent amplification of the scattered waves resulting from the transfer of energy from the pump wave to the scattered wave are dealt with in the preceding section. In order to explore the applicability of the model, a detailed numerical study of threshold condition and gain behaviour of SRS has been made. Efforts are made to find out conditions for achieving larger gain with lower power input and also to find out most suitable material amongst CdS, InSb and GaAs to observr the SRS process. In doing so, we consider the irradiation of n-type doped semiconductor samples (such as CdS, InSb, GaAs) by at room temperature (300K).

The following material parameters are taken as representative values to establish the theoretical formulation and achieve the aim: for CdS; $m_{de}=1.5\times10^{25}m_0m^{-3}$, $m_{dh}=1.2\times10^{25}m_0m^{-3}$, $\rho=4.28\times10^{3}kgm^{-3}$,for GaAs; $m_{de}=4.7\times10^{23}m_0m^{-3}$, $m_{dh}=9.8\times10^{24}m_0m^{-3}$, $\rho=5.28\times10^{3}kgm^{-3}$ for InSb; $m_{de}=4.2\times10^{22}m_0m^{-3}$, $m_{dh}=7.3\times10^{24}m_0m^{-3}$, $\rho=5.28\times10^{3}kgm^{-3}$. These crystals are assumed to be irradiated by a pulsed $10.6\mu m\ CO_2$ laser.

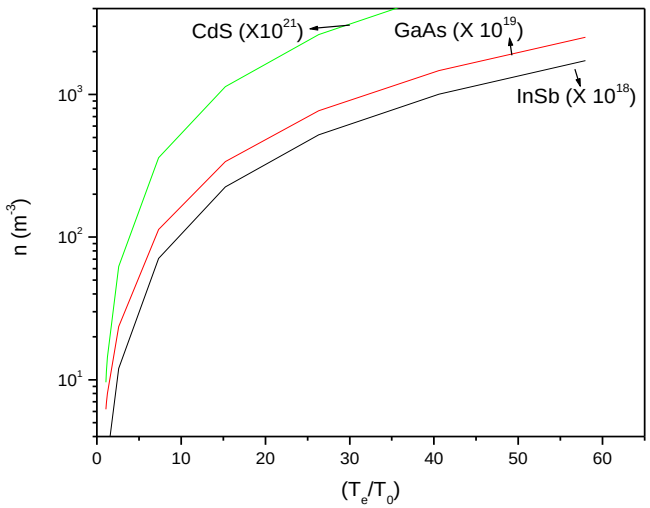

Fig. 1. Variation of maximum plasma density (n) with electron to lattice temperature (T_e/T_0) at $300K$.

Fig.1 depicts the qualitative behaviour of the variation of maximum possible plasma density (n) with carrier temperature (T_e/T_0) for the three direct gap

semiconductors under study. In all the three cases maximum possible plasma density increases with carrier temperature. It is found to be highest in CdS of the order of $(\approx 10^{21}\, m^{-3})$ and lowest in InSb $\left(\approx 10^{18}\, m^{-3}\right)$ crystals. Hence one can conclude from [4] that higher the density of state mass higher is the maximum possible plasma density.

Steady state amplification Characteristics

3.2.1: Variation of Steady state gain with pump amplitude

In this section we spotlight on the numerical analysis of the gain coefficient associated with the SRS with pump field E_0.The pump field should be well above the threshold value (i.e. $|E_0| > |E_{th}|$) to achieve the significant signal. Fig. 2 shows the behaviour of gain coefficient, obtained from eq. (26) with respect to E_0 the pump amplitude. The nature of variation of the gain with pump amplitude is almost similar for all the three semiconductors. Initially for lower pump amplitudes gain increases, leading to a maximum gain point and then decrement is observed on further increasing E_0 value. It is found that to achieve significant amplification in scattered wave, the input pump amplitude should be in between $(1\times 10^{1} Vm^{-1})$ and $(1\times 10^{6} Vm^{-1})$. Most favourable pump amplitude range for maximum gain is from $5\times 10^{2} Vm^{-1}$ to $3\times 10^{5} Vm^{-1}$. Highest gain $(\approx 10^{8}\, m^{-1})$ is obtained for CdS crystal whereas InSb semiconductor gives the lowest gain $(\approx 10^{4}\, m^{-1})$.

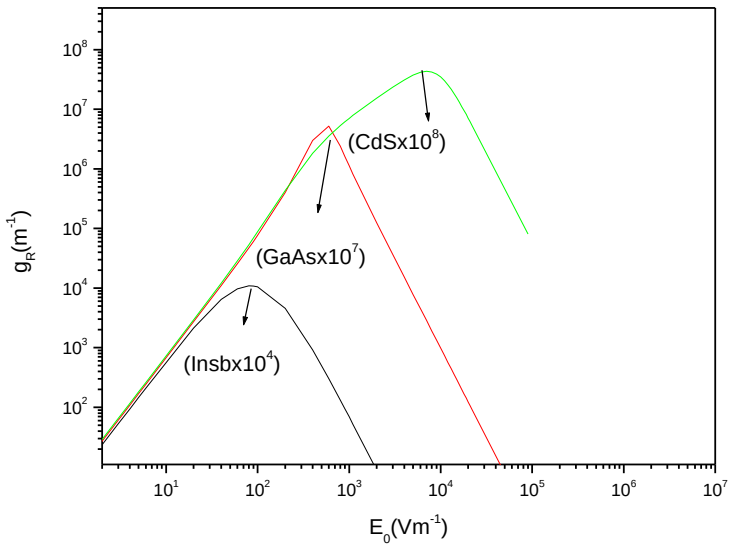

Fig. 2. Variation of Steady state Raman gain coefficient g_R with pump amplitude E_0 at $k = 10^{6}\, m^{-1}$.

3.2.2: Variation of Steady state gain with wave number

Fig. 3 shows the behaviour of gain coefficients as a function of wave number k at constant E_0.The nature of variation of the gains with wave number are almost identical for all the three semiconductors. Initially for lower wave number gain increases with increasing k and then become nearly invariant with k. It is found that to achieve controllable amplification of signal wave, the wave vector values should be in the range $1\times10^{7}\,m^{-1}$ to $8\times10^{8}\,m^{-1}$. Beyond this region, gain becomes nearly independent of wave number. Highest gain $(\approx10^{6}\,m^{-1})$ is obtained for CdS crystal whereas InSb semiconductor gives the lowest gain $\left(\approx10^{3}\,m^{-1}\right)$. Hence from the above discussion one may conclude that highest density of state mass is one of the controllable parameter for steady state gain coefficient for SRS.

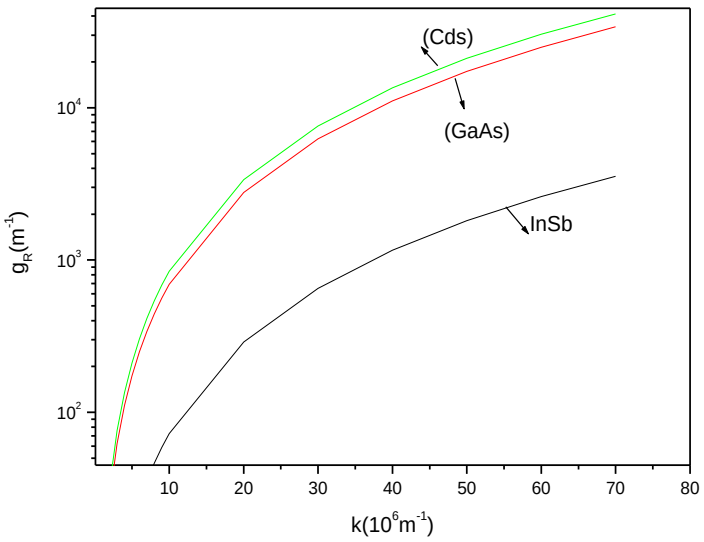

Fig. 3. Variation of Steady state Raman gain (g_R) with wave vector (k) at $E_0 = 10^7 Vm^{-1}$.

3.2: Transient gain characteristics

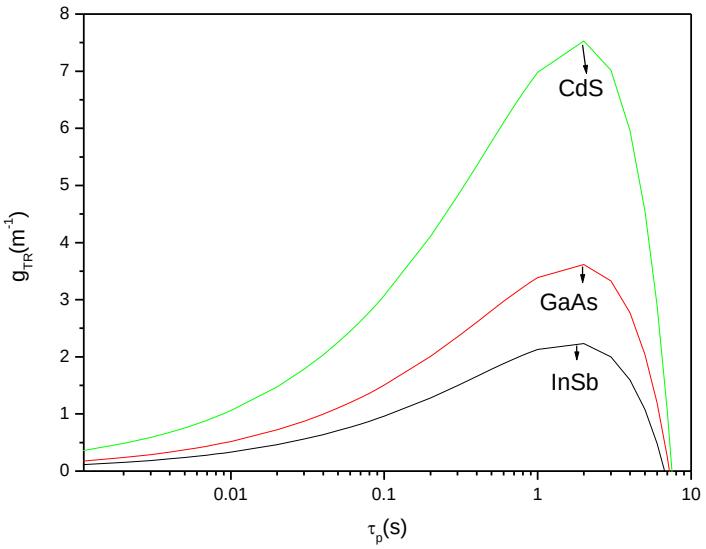

Fig. 4. Variation of Transient Raman gain coefficient g_R with pulse duration τ_p at $I_{in} = 7.646 \times 10^{17} Wm^{-1}$.

Fig. 4 depicts the qualitative behaviour of the transient gain coefficient of signal wave as a function of pulse duration. The cell length is taken in the range of $10^{-4} m - 10^{-6} m$ and pulse duration within the range $10^{-10} \le \tau_p \le 10^{-2} s$ for all the three semiconductor crystals. The density of states mass affects the g_{TR} implicitly and shifts the maximum transient gain point towards higher value of pulse duration. It is clear that behaviour of g_{TR} with respect to τ_p is identical in all the three semiconductor crystals. For a fixed value of I_{in} transient gain factor g_{TR} first increases with τ_p and then attains maximum value which remains constant for certain range of τ_p. If τ_p is further increased g_{TR} diminishes very rapidly and approaches to zero at different pulse durations for the three (CdS, GaAs, InSb) direct gap semiconductor crystals. The maximum transient gain is obtained for CdS and minimum gain is found in InSb. The pulse duration at which transient gain vanishes from InSb for CdS through GaAs.

4. Conclusion

The above analysis supports the fact that the crystals having higher density of state mass are better hosts for the fabrication of the optoelectronic devices. III-V and II-VI direct gap semiconductors are found to be most appropriate host for density dependent stimulated Raman scattering in the allowed range ($\omega_0 \ge \omega_p$) of pump frequency. Such

semiconductors may be used to develop potentially useful optoelectronic devices. CdS is found to be most appropriate host for the SRS process with highest gain coefficient in comparison with InSb and GaAs. But InSb proves its candidature being cost effective material by sacrificing a little towards output.

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