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Photo-induced birefringence in semiconductors compared with optical fibers

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Abstract

We analyze the photon-induced birefringence in semiconductors based on pump-probe setups, within the semiconductor Bloch equations formalism and the Luttinger–Kohn model for the band structure. When the pump and probe pulses are well separated in time, the anisotropic momentum space filling of the photo-excited electrons is the only mechanism causing the induced birefringence. The birefringence ratio is then $B_r = \frac{3}{4}$ for pump and probe having perpendicular vs. parallel linear polarizations. This ratio is $B_r = \frac{1}{6}$ for opposite vs. identical circular polarization. When the pump and probe pulses overlap in time, these birefringence ratios become $B_r = \frac{1}{3}$ for linear polarizations and $B_r = \frac{1}{9}$ in case of circular polarizations. These predictions differ markedly from those for optical fibers.

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1. Introduction

The refractive index of media, like semiconductors and optical fibers, depends on the incident light intensity [1,2]. This optical nonlinearity has been used in various ways of Mach–Zehnder interference modulation switchings, where the refractive index of the media is controlled by a high intensity pump pulse. In case the pump pulse is polarized, either linearly or circularly, the medium will become birefringent due to the excitation. This mechanism has been used extensively to implement the optical polarization rotation switching in both optical fibers [3,4] and semiconductors [5–9].

The induced birefringence in the optical fiber has been well described in the literature [1]. Due to the properties of the virtual excitation in the glass optical fiber, the birefringence relaxation time τ_{bi} and birefringence ratio of the two orthogonal modes B_r, are rather simple. In semiconductors, the light induces also real transitions of the electrons. These electrons contribute to the optical nonlinearity through various processes, like electron valence intraband relaxation, carrier heating, spectral hole burning [10–14]. The description of these phenomena usually focuses on the energy distribution of the carriers, thereby ignoring a possibly created anisotropy in momentum space, which is essentially a source of induced birefringence. That anisotropy is due to the coupling between the orientation of the transition dipole and the momentum of the excited carriers [15–17], as it occurs for instance in $\vec{k} \cdot \vec{p}$ theory of band structure. Intuitively this will lead to birefringence on a time scale of the relaxation times of the carrier populations, which was therefore studied in

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[15–17]. However, this anisotropy turns out to be rather weak, insufficient to explain the observed large phase shift and gain difference between the two orthogonal modes [5–9]. Moreover, due to both *Boltzmann* scattering and "memory effects", the relaxation in the energy space is theoretically faster than that of the momentum orientation relaxation [17]. This, in principle, precludes the observation of the birefringence on an ultrashort time scale.

Here we note that besides the electron populations of the conduction and valence bands also the off-diagonal elements of the density matrix, i.e. the polarization, can contribute to the induced birefringence. We shall argue that this birefringence is quite pronounced. Moreover, it could be of interest for applications in ultrafast switching, as the relaxation times of the polarization densities are expected to be of the order of only a few tens of femtoseconds. We use the well known Semiconductor Bloch equations [18,20,19] to study the dipole polarization nonequilibrium in an optically exited semiconductor, its influence on the response to the polarized electric field, and the resulting birefringence ratio B_r. Both linear and circular polarization birefringence are considered. In the last section, we compare the results obtained for semiconductors with those for optical fibres. In spite of some similarities there are also marked differences.

2. Anisotropic momentum space filling

In this section we first give the analysis of the birefringence caused by anisotropic electron momentum space filling effects due to the excitation with a linearly polarized pump pulse. The experimental situation on which we will focus is that of linearly polarized pump and probe pulses. Their polarization vectors define the x–z plane, with pump polarization $\hat{e}_p = \hat{z}$ and probe polarization $\hat{e}_t = \hat{z}\cos\theta_0 + \hat{x}\sin\theta_0$, so $\theta_0 = 0$ means co-polarization and $\theta_0 = \pi/2$ cross-polarization. These polarization vectors, together with the direction of propagation as the y axis, define the laboratory reference frame. In this frame the Bloch vector $\vec{k}$ of an electron that is involved in an optical transition has polar angles θ and ϕ, as shown in Fig. 1.

The major part of the optical transitions in III–V semiconductors are those between the conduction band and heavy-hole band, of which the wave function is given by [21]

$$\psi_h = -\frac{1}{\sqrt{2}}[(|X\rangle\cos\phi + |Y\rangle\sin\phi)\cos\theta - i(|X\rangle\sin\phi - |Y\rangle\cos\phi) - |Z\rangle\sin\theta] \quad (1)$$

and with the spin parallel to the $\vec{k}$ vector. The components $|X\rangle$, $|Y\rangle$ and $|Z\rangle$ denote the normalized p-like wave functions. By optical transition they are connected to the conduction band $|s\rangle$ state with the same $\vec{k}$ vector and same spin direction.

A pump pulse $\vec{\mathscr{E}}_p(\vec{r},t)$ with polarization along the z direction couples, in dipole approximation $-e\vec{\mathscr{E}}\cdot\vec{r}$, only to the $|Z\rangle$ component of the heavy-hole wave function. Therefore the optical transition probability between the heavy-hole band and the conduction band is for a given direction of $\vec{k}$ proportional to $\sin^2\theta$. In other words, the excited electrons are anisotropically distributed in $\vec{k}$ space, with a $\sin^2\theta$ distribution. This distribution is illustrated in Fig. 2. Electrons excited from the light-hole band have a complementary distribution in $\vec{k}$ space, but are much smaller in number because of the much lower transition state density. They are therefore neglected.

The probe pulse, linearly polarized in the x–z plane at an angle θ_0 relative to the pump pulse, is sensitive to a $|Z\rangle\cos\theta_0 + |X\rangle\sin\theta_0$ component of the hole wave function. The change of the gain (absorption) that it experiences due to the electrons that were excited by the pump is then, after integration over the aforementioned $\sin^2\theta$ distribution of the carriers, found to depend on θ_0 as

$$\delta g_{\text{probe}} \propto N\left(\frac{4}{5}\cos^2\theta_0 + \frac{3}{5}\sin^2\theta_0\right). \quad (2)$$

This induced gain anisotropy Eq. (2) is rather weak, corresponding to a birefringence ratio $B_r = \frac{3}{4}$, assuming an isotropic linewidth enhancement factor. This weak birefringence is difficult to reconcile with the large observed polarization rotations [5–9]. Moreover, the measured relaxation time of the birefringence is in the order of tens of femtoseconds [16], which is rather different from the population distribution relaxation times of a few hundred femtoseconds found in theoretical studies [17].

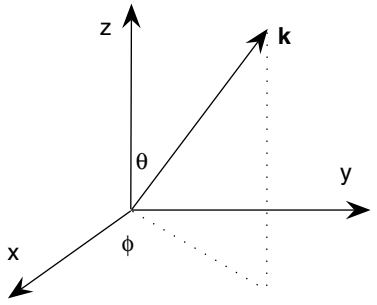

Fig. 1. The orientation of the electron wave vector $\vec{k}$ within a laboratory frame defined by the x–z plane of pump and probe polarizations.

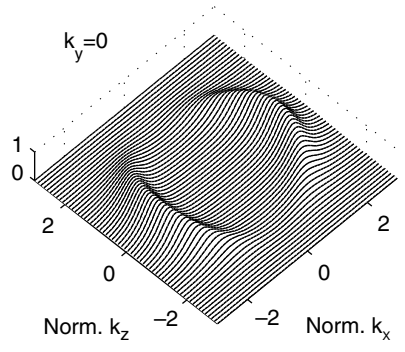

Fig. 2. Anisotropic carrier population density in momentum space excited by a z polarized pulse.