Enhancement of the Nonlinear Acoustoelectric Interaction in a Photoexcited Plasma in a Quantum Well

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Nonlinear interaction of an intense surface acoustic wave (SAW) with a 2D electron–hole plasma generated by light in a semiconductor quantum well near a piezoelectric crystal is investigated. It is shown that, in a strongly nonlinear regime, the acoustoelectric interaction is enhanced because of the accumulation of carriers in the field of an intense SAW. In addition, in a strongly nonlinear regime, the dissipation of the acoustic wave energy increases and the sound velocity decreases. These dependences fundamentally differ from those observed in a unipolar plasma. For high sound intensities, analytical results are obtained.

Surface acoustic waves offer good possibilities for studying of 2D electron systems [1] owing to the piezoelectric interaction. In recent years, the so-called hybrid structures have been developed [2]. A hybrid structure (Fig. 1) consists of a submicron thin film (A_3B_5) with an electron quantum well and a substrate made of a piezoelectric crystal ($LiNbO_3$). Such a system is a quasi-monolithic one, because the film is tightly bound to the piezoelectric crystal by the Van der Waals forces. For a surface acoustic wave (SAW) in a hybrid structure, the electromechanical coupling constant K_{eff}^2 is two orders of magnitude greater than in conventional GaAlAs systems [2]. In a hybrid structure, the regime corresponding to nonlinear acoustoelectric interaction can be obtained in experiments. The experimental studies [3, 4] showed that the electron plasma (which was induced in the cited experiments by a metal gate) could be divided into strips in the piezoelectric field of a SAW.

In this paper, we theoretically investigate the nonlinear acoustoelectric interaction in a hybrid structure in the presence of laser illumination, which generates an electron–hole plasma at room temperature. In hybrid structures, the acoustoelectric interaction is mainly governed by the piezoelectric effect. In the presence of a high-intensity SAW, the mean carrier density increases and the plasma falls into electron and hole strips. As the sound intensity increases, the absorption of the SAW also increases and the SAW velocity shift due to plasma decreases. These dependences are related to a strong recombination nonlinearity. The behavior of a unipolar plasma in an intense acoustic wave is fundamentally different: the absorption of the acoustic wave saturates, and the sound velocity increases [3–6].

In piezoelectric crystals with mobile electrons, the nonlinearity of the acoustoelectric interaction may be of a concentration character [5–7] or may occur as a

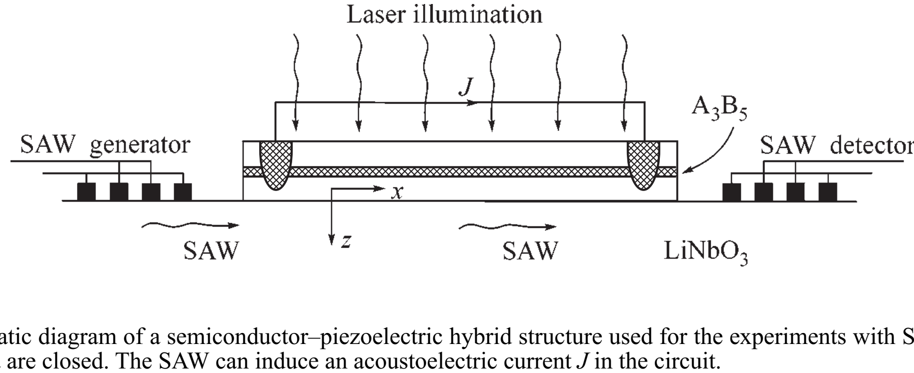

Fig. 1. Schematic diagram of a semiconductor–piezoelectric hybrid structure used for the experiments with SAW. The contacts to the 2D plasma are closed. The SAW can induce an acoustoelectric current J in the circuit.

result of the capture of electrons by traps [8]. In this paper, we consider the combined effect of two mechanisms of nonlinearity, namely, the concentration mechanism and the recombination one. In our model, the recombination nonlinearity occurs in the case of the luminescence of mobile electron–hole pairs in a direct gap semiconductor and in the presence of permanent illumination. The recombination nonlinearity in some respect is similar to the nonlinearity due to traps, because it also leads to a change in the mean density of mobile carriers. However, there are some important differences. We note that, in the cited papers [8], the nonlinearity associated with traps was considered in the absence of illumination and for static traps. In our model, both electrons and holes are mobile. This fact plays a fundamental role, because the mobility leads to a spatial separation of the carriers in the field of an intense SAW.

In this paper, we consider the conventional case of weak acoustoelectric interaction between an intense SAW and the 2D plasma; i.e., we assume that $K_{\mathrm{eff}}^2 \ll 1$. In addition, the sample is considered to be short in the sense that $\Gamma L < 1$ and long in the sense that $L \gg \lambda$. Here, Γ is the absorption coefficient for the SAW, L is the sample length, and λ is the wavelength. When the condition $\Gamma L < 1$ is valid, the contributions of higher harmonics of the SAW play a negligible role. To solve the problem in the framework of the aforementioned approximations, it is sufficient to determine the nonlinear response of the plasma to the piezoelectric potential of a monochromatic SAW. In the nonlinear regime, the SAW velocity shift due to the 2D plasma, δv_s, and the sound absorption coefficient Γ are determined by the expressions [4]

$$\frac{\delta v_s}{v_s^0} = \frac{\langle j\Phi^{SAW}\rangle}{2I}, \quad \Gamma = \frac{\langle jE_x^{SAW}\rangle}{I}, \tag{1}$$

where $\langle\ldots\rangle$ denotes spatial averaging, v_s^0 is the sound velocity in the absence of plasma, I is the SAW intensity, Φ^{SAW} is the piezoelectric potential of the SAW, E_x^{SAW} is the electric field induced by the SAW in the plane of the 2D plasma, and j is the current of the charge carriers. A SAW that propagates in the x direction generates an electric field $E_x^{SAW} = E_0\cos(qx - \omega t)$ and the corresponding piezoelectric potential $\Phi^{SAW} = -\Phi_0\sin(qx - \omega t)$, where $\Phi_0 = E_0/q$, t is time, q is the wave vector, and $\omega = v_s^0 q$. Since we assumed that $K_{\mathrm{eff}}^2 \ll 1$, we have $\delta v_s \ll v_s^0$ and $\Gamma \ll q$.

The current $j = j_n + j_p$ is determined from the system of hydrodynamic equations

$$e\frac{\partial n}{\partial t} + \frac{\partial j_n}{\partial x} = (G - Cnp)e,$$
$$|e|\frac{\partial p}{\partial t} + \frac{\partial j_p}{\partial x} = (G - Cnp)|e|, \tag{2}$$
$$j_n = \sigma_n E_x - eD_n\frac{\partial n}{\partial x}, \quad j_p = \sigma_p E_x - |e|D_p\frac{\partial p}{\partial x},$$

where n and p are the 2D electron and hole densities, respectively; j_n and j_p are the electron and hole currents; $e = -|e|$ is the electron charge; G is the rate of carrier generation by the laser light; C is the recombination constant; σ_n and σ_p are the electron and hole conductivities; and D_n and D_p are the electron and hole diffusion coefficients. Here, the following relations are valid: $\sigma_n = |e|\mu_n n$ and $\sigma_p = |e|\mu_p p$, where μ_n and μ_p are the electron and hole mobilities and $D_n/\mu_n = D_p/\mu_p = KT/|e|$, where T is temperature.

The electric field consists of two components: $\mathbf{E} = \mathbf{E}^{SAW} + \mathbf{E}^{ind}$. The field $\mathbf{E}^{ind}$ is induced by the plasma and determined from the Maxwell equation

$$\mathrm{div}\mathbf{D}^{ind} = 4\pi(en + |e|p)\delta(z), \tag{3}$$

where $\mathbf{D}^{ind} = \hat{\epsilon}(z)\mathbf{E}^{ind}$ and z is normal to the system. The 2D plasma in the quantum well is modeled by the δ-function.

We will seek a steady-state periodic solution to the system of Eqs. (2) and (3) as a function of the coordinate $x' = x - v_s^0 t$. We note that, in our system with shorted contacts to the 2D plasma, we have $\langle E_x\rangle = 0$ (see Fig. 1) and the solution can be found in the form of a periodic function of x'. To determine the numerical solution, it is convenient to expand all functions in Eqs. (2) and (3) in Fourier series in the x' coordinate. Neglecting the thickness of the film, we obtain $\Phi_m^{ind} = 2\pi(en_m + |e|p_m)/\epsilon_{\mathrm{eff}}|q_m|$, where Φ_m^{ind} represents the Fourier coefficients; $m = 0, \pm1, \pm2, \ldots, q_m = qm$; and $\epsilon_{\mathrm{eff}} = (\epsilon + 1)/2$. Here, $\epsilon = \sqrt{\epsilon_{xx}^T\epsilon_{zz}^T}$ is the mean permittivity of $LiNbO_3$. The SAW intensity and the amplitude of the potential Φ_0 are related by the formula $K_{\mathrm{eff}}^2 = (\sigma_m q/2)(\Phi_0^2/I)$, where $\sigma_m = \epsilon_{\mathrm{eff}} v_s^0/2\pi$ [4]. The coefficient K_{eff}^2 was calculated earlier in the framework of the linear theory [2, 9].

The system of Eqs. (2) and (3) was solved numerically. The calculations were performed by using about 100 spatial harmonics and the following typical parameters of the system at $T = 300$ K: $\mu_e = 2 \times 10^3$ cm^2/(V s), $\mu_e/\mu_p = 6$, $v_s^0 = 3.9 \times 10^5$ cm/s, $\lambda = 60$ μm, $K_{\mathrm{eff}}^2 =$

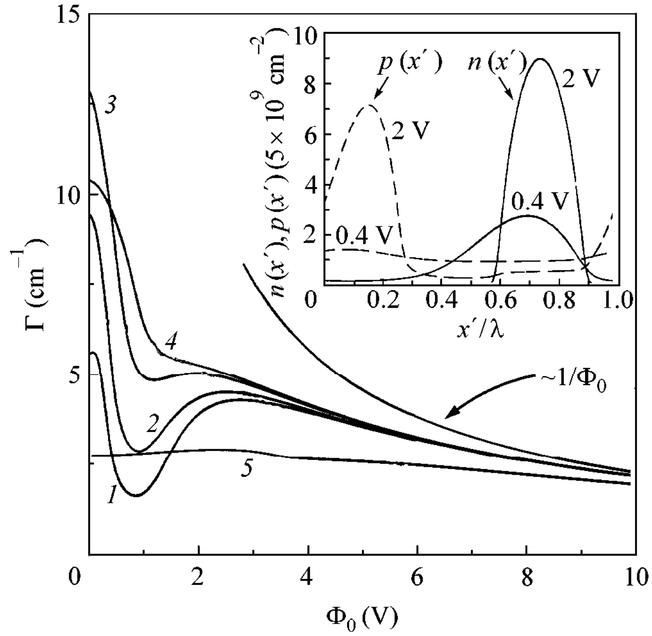

Fig. 2. Dependence of the sound absorption coefficient on the amplitude of the SAW potential. Curves *1–5* correspond to different intensities of optical pumping: $n_{s0} = 10^9$, 2×10^9, 5×10^9, 10^{10}, and 5×10^{10} cm^{-2}. The inset shows the profiles of the electron and hole concentrations for $n_{s0} = 5 \times 10^9$ cm^{-2} and for two different amplitudes $\Phi_0 = 0.4$ and 2 V.

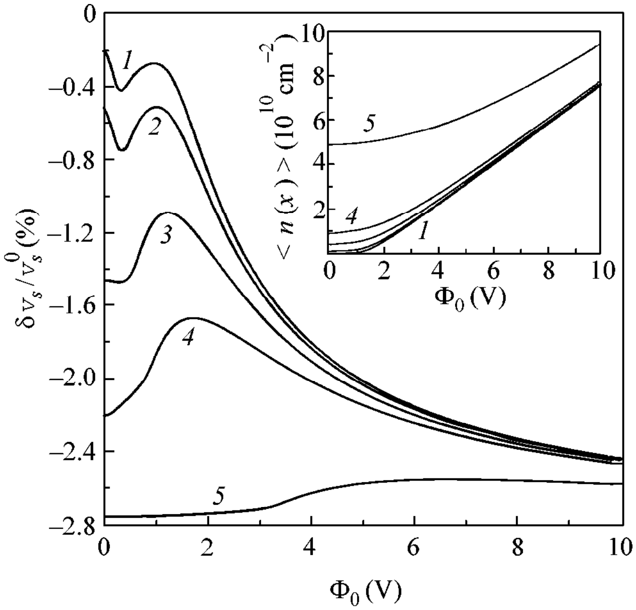

Fig. 3. Dependence of the SAW velocity shift on the SAW amplitude. The curves correspond to those in Fig. 2. The inset shows the dependence of the mean concentration on the SAW amplitude for different values of n_{s0}.

0.056, and $\epsilon = 50$. The recombination constant for nonequilibrium electrons and holes was estimated by a formula similar to Eq. (13) from [10]: $C = 4 \times 10^{-4}$ cm^2/s, which corresponds to the radiative recombination time $\tau_0^{rad} = 1/Cn_{s0} = 250$ ns for the electron and hole densities $n_{s0} = p_{s0} = 10^{10}$ cm^{-2}. At $\Phi_0 = 0$, the steady-state carrier concentrations were $n_{s0} = p_{s0} = \sqrt{G/C}$. The concentration n_{s0} was selected within the interval 10^9–10^{11} cm^{-2}, which is typical of optical experiments [10]. The corresponding luminous fluxes were within 10^{-2}–10^2 W/cm^2. The calculations were performed for the amplitudes of the potential $\Phi_0 \sim 1$–10 V. The amplitudes $\Phi_0 \sim 3$ V are typical of the experiments on hybrid structures with metallic gates [3, 4]. In these experiments, the input powers of SAW generation varied from –10 to 30 dB m [3, 4]. In this paper, we consider a gateless structure, for which the coefficient K^2_{eff} should be greater.

The inset in Fig. 2 shows the calculated functions $n(x')$ and $p(x')$ for different values of Φ_0. One can see that, as Φ_0 increases, the plasma becomes separated into adjacent electron and hole strips. Electrons and holes tend to screen the potential $\Phi^{SAW}(x')$ at all x', but, simultaneously, the functions $n(x')$ and $p(x')$ slightly overlap. Only this type of carrier distribution can exist in a steady-state regime with a strong recombination nonlinearity, because the appearance of a noticeable gap between the strips should lead to a nonstationary accumulation of carriers. The absorption coefficient Γ (Fig. 2) mainly decreases with increasing Φ_0. The nonmonotonic behavior of Γ at $\Phi_0 \sim 1$–2 V correlates with the instant of noticeable separation of the plasma into strips and with the onset of the increase in the mean density $N_0 = \langle n(x')\rangle = \langle p(x')\rangle$ (the inset in Fig. 3). The increase in N_0 is a natural consequence of the spatial separation of electrons and holes and of the suppression of their recombination. When $\Phi_0 \longrightarrow \infty$, the sound absorption coefficient behaves as $\Gamma \propto 1/\Phi_0$ and the absorbed energy is $Q = I\Gamma \propto \Phi_0$. The increase in Q with growing sound intensity is related to the increase in the mean carrier density N_0 with $\Phi_0 \longrightarrow \infty$ (the inset in Fig. 3). Below, we will analytically derive the asymptotic expression for $Q(\Phi_0)$.

The calculations were also performed for the mean currents induced by the SAW. With increasing Φ_0, the current $\langle j_n(x') + j_p(x')\rangle$ first increased in magnitude and then decreased, because the carriers were trapped by the SAW. The mean velocities of electrons and holes tend to v_s^0 with $\Phi_0 \longrightarrow \infty$.

The velocity shift δv_s (Φ_0) shown in Fig. 3 exhibits a maximum at $\Phi_0 \sim 1$–2 V. This maximum is also related to the onset of plasma separation into strips. When $\Phi_0 \longrightarrow \infty$, the velocity shift δv_s decreases, because the plasma becomes more and more dense and strongly screens the field of the SAW.

From the curves shown in the inset in Fig. 3, one can see that the transition to the nonlinear regime of carrier accumulation depends on the initial density $n_{s0} = p_{s0}$. The characteristic value of Φ_0 at which the density N_0 begins to grow increases with n_{s0}. The order of magnitude of this characteristic amplitude $\Phi_0^{non-lin}$ can be estimated in terms of the linear theory. Using the condition $\delta\rho_s \sim |e|n_{s0}$, where $\delta\rho_s$ is the perturbation of the 2D charge, we obtain $\Phi_0^{non-lin} \sim |e|n_{s0}\lambda/\epsilon_{\text{eff}}$. This estimate is valid on condition that $\sigma_0 > \sigma_m$, where $\sigma_0 = |e|(\mu_n + \mu_p)n_{s0}$. For $n_{s0} \sim 10^{10}$ cm^{-2}, we obtain $\Phi_0^{non-lin} \sim 2$V.

The asymptotic behavior of Γ and δv_s for $\Phi_0 \longrightarrow \infty$ can be obtained analytically from the following consideration. When $\Phi_0 \longrightarrow \infty$, the plasma screens the SAW field E_x^{SAW} at practically all points, and the density of 2D carriers can be determined from the simple integral equation $E_x = E_x^{ind}(x') + E_x^{SAW}(x') = 0$. In this case, the electrons screen half of the wave in the region where $e\Phi^{SAW} < 0$, and the holes screen half of the wave in the region where $e\Phi^{SAW} > 0$. From the condition $E_x = 0$, we obtain the asymptotic expression $N_0 \longrightarrow q\Phi_0\epsilon_{\text{eff}}/2\pi^2|e|$, which adequately describes the numerical data given in the inset in Fig. 3. According to the Drude model, the local absorption of the SAW by the plasma has the form

$$Q_{loc} = \frac{v_n^2 m_e^*}{\tau_n} n + \frac{v_p^2 m_h^*}{\tau_p} p,$$

where v_n and v_p are the velocities of electrons and holes, m_e^* and m_h^* are the electron and hole effective masses, and τ_n and τ_p are the corresponding relaxation times. Since the electrons and holes are almost totally trapped by the SAW and spatially separated, we have $v_n \simeq v_s^0$ for $\lambda/2 < x < \lambda$ and $v_p \simeq v_s^0$ for $0 < x < \lambda/2$. The total absorption now takes the form

$$Q = \langle Q_{loc} \rangle \simeq (v_s^0)^2 N_0 \left(\frac{m_e^*}{\tau_n} + \frac{m_h^*}{\tau_p}\right)$$
$$= (v_s^0)^2 N_0 |e| \left(\frac{1}{\mu_n} + \frac{1}{\mu_p}\right)$$
$$= (v_s^0)^2 \left(\frac{1}{\mu_n} + \frac{1}{\mu_p}\right) q\Phi_0\epsilon_{\text{eff}}/2\pi^2.$$

Thus, we derived the asymptotic expression $Q \propto N_0 \propto \Phi_0$ for $\Phi_0 \longrightarrow \infty$. The small parameters corresponding to $\Phi_0 \longrightarrow \infty$ are $v_s^0/\mu_n E_0 \ll 1$ and $v_s^0/\mu_p E_0 \ll 1$. To determine δv_s for $\Phi_0 \longrightarrow \infty$, we used similar small parameters and obtained the asymptotic expression $\delta v_s/v_s^0 \longrightarrow -K_{\text{eff}}^2/2$.

In the case of a unipolar plasma, which is induced, e.g., by the gate, the quantity $\langle n \rangle$ will evidently be constant for any Φ_0. When $\Phi_0 \longrightarrow \infty$, electrons form narrow strips and, on the whole, the screening of the SAW field is reduced. In the limit $\Phi_0 \longrightarrow \infty$, we have $Q \longrightarrow Q_{\max} = \langle n \rangle (v_s^0)^2 m_e^*/\tau_e$ and $\delta v_s \sim 1/\Phi_0 \longrightarrow 0$ [4–6].

The experimental studies of acoustoelectric interaction in 3D crystals in the presence of illumination were performed in the regime corresponding to hopping photoconduction and weak nonlinearity (see, e.g., [11]). The acoustoelectric interaction between the SAW and the photoexcited plasma in GaAs gave rise to a longitudinal voltage between the gates [12]. A strongly nonlinear acoustoelectric interaction in 2D systems was experimentally studied at $T = 4$ K in the exciton ionization regime [13, 14]. For example, a delay line for photons was realized in [13]: the incident photons were converted to electrons and holes and then carried along the sample by the SAW; at the final stage, the carriers generated secondary photons. At $T = 300$ K, an intense SAW was used to study a strongly nonuniform plasma excited in a quantum well by a laser beam [15]. One more class of experiments is related to the diffraction of light by volume and surface acoustic waves [16].

The predicted nonlinear acoustoelectric interaction applies to uniform illumination in an infinite sample. In a finite sample, the density accumulated in the potential wells of the SAW saturates when $\Phi_0 \longrightarrow \infty$. In our system, an intense SAW first separates the electrons and holes and then carries them along the sample. At the final stage, the carriers recombine as soon as they reach the right contact. In other words, an intense SAW "collects" the carriers over the whole sample length. The maximum density accumulated by SAW is equal to $N_{\max} = GL/v_s^0 = (L/v_s^0)n_0^2 C = n_0(t_0/\tau_0^{rad})$, where $t_0 = L/v_s^0$ is the time of SAW propagation over the sample length. Thus, in a finite system, each curve shown in the inset in Fig. 3 is saturated at its own level $N_{\max}(n_0)$. The described enhancement of the acoustoelectric interaction in the nonlinear regime manifests itself when $N_{\max} > n_0$. The latter inequality is equivalent to $t_0 > \tau_0^{rad}$. These inequalities hold for the typical parameters of our problem. For example, when $n_0 = 10^{10}$ cm^{-2} and $L = 2$ mm, we obtain $t_0/\tau_0^{rad} \sim 2$.

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