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ACOUSTOELECTRIC EFFECT DUE TO A PIZOELECTRICALLY ACTIVE WAVE IN MANY-VALLEY SEMICONDUCTORS OF THE TYPE n–GaP AND n-AlSb

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The theory of acoustoelectronic effects in many-valley piezoelectrics should differ from the "single-valley" case. The objective of this paper is to determine the degree of this difference. Accordingly, we calculate the acoustoelectric current (j^{ae}) and field (E^{ae}) as well as the electron absorption coefficient (Γ_e) and the variation of the acoustic wave velocity (∇v_s) for a piezoelectrically active acoustic wave with $q \,||<110>$ and polarization vector $e^{\circ}||<110>$ in semiconductors of the n-GaP type, and we investigate the influence of $\upsilon_{\alpha\alpha'}$ on j^{ae}, E^{ae}, Γ_e, ∇v_s. .

KEYWORDS : Many-valley semiconductors, acousto-electronic coupling force, polarization vector, piezoelectric polarization.

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INTRODUCTION

Acoustoelectric effects in a many–valley semiconductors have been investigated heretofore only in the case of nonpiezoelectric material [1-6], in which the interaction of an acoustic wave with current carriers is realized through the deformation potential. The results are also applicable to piezoelectrics, but in the case of "piezoelectrically inactive" wave [7]. However, the analogous effects associated with pizeoelectrically active waves can be many times stronger [7, 8]. Several many-valley pizeoelectrics are known and are finding important engineering application at the present time, for example GaP, GaAs, CdS, AlSb, AlAs and AlP (n-type), in which the lowest conduction band has the structure of the n-Si conduction band [9-15]. It might seem at first glance that the theory developed for "single valley" semiconductor would be applicable to these crystals. Indeed, the piezoelectric potential, unlike the deformation potential, displaces all valleys identically, so that the intervalley transitions, so it would seem, should not influence acoustoelectronic effects. That being the case, the net effect of mutual influence of the valleys would merely be their contribution to the total screening. If the latter is weak, the crystal should behave as a set of independent valleys.

However, such is not the case. As a matter of fact, the action of an acoustic wave on electrons of the α-th valley is determined, not by the potential, but by the acoustoelectronic

coupling force ($-\nabla\varphi$), which is proportional to the acoustic wave vector q. Due to the anisotropy of the effective mass, the different orientations of q relative to the principal axes of different valleys elicit unequal accelerations of the electrons in them, even when the force is identical. This situation causes unequal heating of the electron gas, an effect that is manifested in differences between the symmetrical parts of the nonequilibrium distribution functions in different valleys (another factor responsible for this effect is scattering anisotropy, *i.e.*, the tensor nature of the momentum relaxation time). As a result, the "partial" (for the carrier energy ε) dynamic equilibrium between valleys is upset, inducing a redistribution of electrons in them, if the probability $\upsilon_{\alpha\alpha'}$ of intervalley Umklapp events is much lower than the acoustic frequency ω, *i.e.,* if such events are rare, then the heating differences are a maximum. If, on the other hand, $\upsilon_{\alpha\alpha'} >> \omega.$, then these differences are annihilated by rapid intervalley mixing. But this means that "piezoelectrically active" acoustoelectronic effects will depend on $\upsilon_{\alpha\alpha'}$. However, unlike the case of deformation potential interaction [1-10], this dependence is a purely anisotropic effect.

THEORY AND CALCULATION

The method of calculation generalizes the method of Gantsevich and Gurevich [4], differing from the latter in that : (1) The electric potential φ generated by the piezoelectric polarization and space charge is included in the kinetic equation; (2) the neutrality condition is replaced by the Poisson equation for a piezoelectric semiconductor, permitting an arbitrary degree of screening to be taken into account; (3) the number (N_v) and positions of the valleys as well as $q^\circ \equiv q^\circ q^{-1}$ and a° are assumed to be arbitrary in the initial equations, and the problem is reduced to a coupled system of N_v kinetic equations and the Poisson equation; (4) the intervalley elastic collision operator is approximated by the tensor $\hat{\tau}_P^\alpha$ and $\hat{\tau}_P^\alpha$ and $\upsilon_{\alpha\alpha'}$ are assumed to have an arbitrary dependence on the carrier energy ε (however intervalley transitions are assumed to take place at impurities, as is valid for sufficiently low temperatures). The calculations are confined to the domain.

$$\tau_p << \tau_i << \tau_e;\ ql_p < 1, \qquad \ldots(1)$$

where τ_p, τ_e, τ_i are the momentum, energy, and intervalley relaxation times, respectively, and l_p, is the "momentum" mean free path of the carriers. The parameters $\omega\tau_i$ and $\omega\tau_e$ can be arbitrary. (The inequality $\tau_i << \tau_e$ is substantiated in [7] for *n*-Ge.) Nondegenerate statistics are postulated. Expanding in terms of the small parameter $\Psi = e\tilde{\varphi}/\kappa_0 T$, where $\tilde{\varphi}$ is the sum of the electric and deformation potentials, we obtain a system of N_v equations for the complex amplitudes of the isotropic nonequilibrium parts of the distribution functions (f_ε^α).

$$\sum_{\alpha'}^{N_v} p_{\alpha z'} f_{\alpha'} = i\varphi_\alpha f_0;\ f_\alpha \equiv f_z^\alpha + \tilde{\varphi}_\alpha f_0;\ \tilde{\varphi}_\alpha = (e\varphi + iqU_0 C_{eq}^\alpha)(\kappa_0 T)^{-1};$$

$$p_{\alpha\alpha'} \equiv \frac{\upsilon_{\alpha\alpha'}}{\omega} = \frac{V(\overline{2m})^{\frac{3}{2}}}{4\pi^2\hbar^3\omega}\sqrt{e}W_{\alpha\alpha'}(\varepsilon,\varepsilon),(\alpha' \neq \alpha);$$

$$p_{\alpha\alpha} = i - \gamma_\alpha - \sum_{\alpha' \neq \alpha} p_{\alpha\alpha'};\ \gamma_\alpha = q v_s^{-1} D_{qq}^{\alpha}(\varepsilon);\ D_{qq}^{\alpha}(\varepsilon) = << (q^0 \cdot v)(q^0 \cdot \hat{\tau}_\rho^{\alpha} v) >>;$$

$$W_{\alpha\alpha'}(\varepsilon, \varepsilon') = \frac{1}{(4\pi)^2} \int d\Omega \int d\Omega'\, W_{\kappa\kappa'}^{\alpha\alpha'}. \qquad \ldots (2)$$

Here f_0 is the equilibrium distribution in one valley, φ is the complex amplitude of the electric potentials, U_0 is the initial displacement amplitude in the acoustic wave, C_{eq}^{α} is the convolution of the deformation-potential tensor in the α-th valley with the vectors e^{o} and q^{o}, $\bar{m}$ is the effective mass of the density of states in one valley, the brackets <<…>> denote averaging over the isoenergetic surface, such averaging in the α-th and α′-th valleys is implicit in $W_{\alpha\alpha'}, v_s$ is the acoustic wave velocity $W_{\kappa\kappa'}^{\alpha\alpha'}$, $\delta(\varepsilon - \varepsilon')$ is the probability of the inter-valley transition $(\kappa, \alpha) \to (\kappa', \alpha')$, and V is the volume of the crystal.

On the basis of the second-order kinetic equations in ψ we express $j^{\alpha e}$ in terms of $f_{\varepsilon}^{\circ\alpha}$ and ψ_α:

$$j^{\alpha e} = \frac{n_0 q \kappa_0 T}{2} \langle \overline{f_0^{-1} \operatorname{Re}(i \tilde{\psi}_\alpha^* f_\varepsilon^{\alpha}) \hat{Z}(q^0 \mu_s^{\perp} + q_\alpha^0 \Delta\mu_s)} \rangle$$

$$\hat{Z}(...) = x^{-\frac{3}{2}} \frac{d}{dx}(x^{\frac{3}{2}}...);\ x = \frac{\varepsilon}{\kappa_0 T};\ \Delta\mu_\varepsilon = \mu_\varepsilon^{\parallel} + \mu_\varepsilon^{\perp};$$

$$\mu_\varepsilon^{\parallel,\perp} = e(m_{\parallel,\perp})^{-1} \cdot \tau_{\parallel,\perp}(\varepsilon). \qquad \ldots (3)$$

Here q_α^0 is the component of the vector q^0 along the axis of rotation of the α-th valley (energy ellipsoid), $(m_{\parallel,\perp})^{-1}$ and $\tau_{\parallel,\perp}$ are the principal components of the reciprocal effective mass tensor and the tensor $\hat{\tau}_P^{\alpha}$, and the brackets <…> denote the usual average of kinetic theory. The overbar denotes arithmetic averaging over all valleys.

The quantities $f_{\varepsilon}^{\circ\alpha}$ and ψ_α are determined from the system (2) and the Poisson equation for a piezoelectric semiconductor; Γ_e and ΔV_s can be determined from the equation of motion of an elastic piezoelectric many-valley medium and the Poisson equation by expressing them in terms of $f_{\varepsilon}^{\circ\alpha}$ and the corresponding coupling constants (deformation-potential and piezoelectric). We now note that for a wave with $q||<110>$, $e^{\circ}||<110>$ and polarization vector $e^{\circ}||<110>$ we have $C_{eq}^{\alpha} = 0$ for all α. Thus, using the well-known Herring equation for the deformation-potential tensor,

$$C_{eq}^{\alpha} = \xi_d \delta_{ij} + \xi_u K_i^{\alpha} K_j^{\alpha},$$

where K_i^{α} denotes the projections of the unit vector directed from the centre of the Brillouin zone to the centre of the α-th valley and ξ_d, ξ_u are the Herring constants, we obtained.

$$C_{eq}^{\alpha} = C_{ij}^{\alpha} e_i^0 q_j^0 = \frac{1}{\sqrt{2}}(C_{31}^{\alpha} + C_{32}^{\alpha}) = \frac{\xi_u}{\sqrt{2}} K_3^{\alpha}(K_1^{\alpha} + K_2^{\alpha}). \quad \ldots(4a)$$

Expression (4a) is clearly equal to zero for valleys centered on the coordinate axes. Thus, the given wave is "deformation-potential inactive", but of course it is piezoelectrically active, because the corresponding coupling constant.

$$C_p = \frac{4\pi e}{qx_0}\beta_{qeq} = \frac{4\pi e}{qx_0}\beta_{123} \neq 0 \quad \ldots (4b)$$

(where β_{qeq} is the convolution of the piezoelectric tensor with the vectors q^0 and e^0, κ_0 is the permittivity of the medium and e is the carrier charge). The calculations yield the following expressions for Γ_e, ΔV_s, $j^{\alpha e}$, $E^{\alpha e}$:

$$\Gamma_e = \frac{\Gamma_0}{\Delta_0} S_r; \quad \frac{\Delta v_s}{v} = \frac{\Gamma_0}{2q\Delta_0}[S_i\eta(S_r^2 + S_i^2)];$$

$$j^{ae} = -\sigma E^{ae} = \frac{J\Gamma_0}{v_s\Delta_0}\langle\overline{y_a\hat{Z}\tilde{\mu}_{\alpha}}\rangle; \Gamma_0 = \frac{qn_0C_p^2}{\rho v_s^2 k_0 T};$$

$$\Delta_0 = \left(1-\eta S_i\right)^2 + \left(\eta S_r\right)^2; \eta = \left(qr_D\right)^{-2;}\ S_r = \langle\frac{3}{2x}\overline{y_{\alpha}}\rangle;$$

$$S_i = \langle\frac{3}{2x\Delta_1}\overline{b_{\alpha}}\rangle - 1; \Delta_1 = B^2 + G^2; B = 6p_{13}\overline{b_a};$$

$$G = 1 - 6p_{13}\overline{g_{\alpha}}; g_{\alpha} = \tilde{\gamma}b_{\alpha}; \tilde{\gamma}_{\alpha} = \gamma_{\alpha} + 6p_{13};$$

$$b_a = (1+\tilde{\gamma}_{\alpha}^2)^{-1}; y_{\alpha} = \Delta_1^{-1}(g_{\alpha}G - b_{\alpha}B);$$

$$\tilde{\mu}_{1,2} = \frac{1}{2}(\mu_{\varepsilon}^1 + \mu_{\varepsilon}^{\perp}); \tilde{\mu}_3 = \mu_{\varepsilon}^{\perp}; \gamma_{1,4} \equiv \gamma_x = \gamma_{2,5} \equiv \gamma_y \neq \gamma_{3,6} = \gamma_z. \quad \ldots (5)$$

Here J is the acoustic energy flux density vector, ρ is the density of the medium, n_0 is the equilibrium number density of carriers in the crystal, r_D is the screening radius, and $p_{1,2}$ is the probability of transitions between valleys situated on different coordinate axes and is, as will become apparent, a quantity on which the acoustoelectronic effects depend. This dependence vanishes in the isotropic case, as can be verified on the basis of (5).

In the case of ultimately slow intervalley relaxation ($\omega\tau_i >> 1$; $q^2D\tau_i >> 1$, where D is the average diffusivity) and weak screening ($\eta << 1$), we infer from (5) the exact results for j^{ae} and Γ_e from the corresponding theory for single-valley semiconductors, summed over all valleys. The comparison in this case must be made in the "mesoultrasonic" range ($\omega\tau_e >> 1$; $ql_p < 1$), because the energy relaxation in (5) is neglected in accordance with (1).

In the opposite case of ultimately fast intervalley relaxation ($\omega\tau_i << 1$; $q^2D\tau_i << 1$) and arbitrary screening, we infer from (5) the results of in which the partial mobility tensor $\hat{\mu}_{\varepsilon} \equiv \mu(x)$ is replaced by the arithmetic average (over all valleys) of the corresponding tensors for the individual valleys. The reason for this is that rapid inter-valley mixing, in effect, creates a single valley with the average mobility.

A comparison of these limiting cases enables us to determine the significance of the dependence of j^{ae}, Γ_e etc., on $\upsilon_{\alpha\alpha'}$ in the investigated crystals. An investigation has shown that this dependence is essentially very weak. A variation of $\upsilon_{\alpha\alpha'}$ over the entire interval between the indicated limiting cases does not change j^{ae}, Γ_e, Δv_s, E^{ae} by more than a few percent. The reason for this result is the following. In *n*-GaP AlSb, AlAs, and AlP, the isoenergetic ellipsoids are greatly elongated along the axes of rotation ($m_{\square} >> m_{\perp}$). The vector q is directed relative to the ellipsoids on the x, y, and z axes in such a way that the effective partial mobilities are, respectively,

$$\tilde{\mu}_{x,y} \approx \frac{1}{2}\mu_{\varepsilon}^{\perp};\ \tilde{\mu}_z \approx \mu_{\varepsilon}^{\perp}(\mu_{\varepsilon}^{\square} << \mu_{\varepsilon}^{\perp}).$$

Thus, for an arbitrarily strong anisotropy of the partial mobility in one valley ($\hat{\mu}_{\varepsilon}^{\alpha}$), the anisotropy of $\tilde{\mu}_{\alpha}$ turns out to be comparatively slight. Consequently, the differences in the "heating" of the valleys will also be slight. Clearly, another significant consideration is the fact that the number of electrons with effective mobility $(1/2)\mu_{\varepsilon\perp}$ is twice that of electrons with $\mu_{\varepsilon\perp}$. Therefore, the pair of valleys $\alpha = x, y$ exists in effective dynamic equilibrium with, the valley $\alpha = z$. It is interesting to note that if the ellipsoids were highly prolate ($m_{\square} << m_{\perp}$), the effective mobilities $\tilde{\mu}_{x,y}$ and $\tilde{\mu}_z$ would differ greatly, since $\tilde{\mu}_{x,y} = (1/2)\mu_{\varepsilon\square}; \tilde{\mu}_z = \mu_{\varepsilon\perp}$, whereupon j^{ae} and Γ_e would depend strongly on $\upsilon_{\alpha\alpha'}$.

The inferred virtual independence of piezoelectrically active acoustoelectronic effects on the probability of intervalley transitions in crystals of the *n*-GaP type is of heuristic significance, because it extends to a broader range of conditions than in the present article. It is valid under the strong influence of carrier traps, electron-electron collisions, various nonlinearities, and in the presence of an external electric field aligned with q (*i.e.*, for the theory of sound amplification), because the cause of the small influence of $\upsilon_{\alpha\alpha'}$ remains in effect. In all of these cases, the allowance for $\upsilon_{\alpha\alpha'}$ would make the theory inordinately cumbersome.

Conclusion

We note in conclusion the existence of another piezoelectrically active wave in cubic semiconductors: $e^{\circ} = q^{\circ}||<111>$. Because of its complete symmetry with respect to all valleys of the type <100> the influence of $\upsilon_{\alpha\alpha'}$ on j^{ae} and Γ_e vanishes.

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