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Acoustic phonon modulation of the electron response in tunnel-coupled quantum wells

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Modulation of the electron properties of tunnel-coupled double quantum wells by acoustic phonon pumping is considered. Phonon-induced voltage (or induced dipole moment) and phonon-drag effect are calculated for the case of high-energy phonon excitation. Linear electron response, with respect to the phonon-pumping intensity, is studied taking into account intrawell and interwell scattering processes due to interface roughness. From numerical estimations we obtained a phonon-induced drop in the potential on the order of 0.1–0.01 meV and a characteristic drag velocity on the order of 100 cm/s for a nonequilibrium phonon temperature equal to 20 K. The possible applications of such structures are noted. [S0163-1829(97)03747-8]

I. INTRODUCTION

Electron response in tunnel-coupled double quantum wells (DQW's) on different types of perturbations has been intensively investigated recently. It was shown that electron transition in DQW's qualitatively differ from the single QW case due to the possibility of electron tunneling between coupled QW's.[1,2] Such an effect is pronounced in peculiarities of the longitudinal transport[3,4] and of the magnetotransport.[5,6] Modification of the electron states and selection rules in DQW's due to tunneling has been also observed in the interband[7,8] and intersubband[9,10] optical spectra. Transition between tunnel-coupled ground states in DQW's under THz (far-infrared) radiation were analyzed.[11,12] A typical ground-state splitting in DQW's is on the order of several meV so that electron transitions between them may be excited not only by THz radiation but also by acoustic phonon pumping in the meV energy range. Such mechanisms of the intersubband transition in DQW's have been discussed[13] in connection with acoustic phonon emission by nonequilibrium electrons. It was noted also that the different types of transitions have taken place for modes propagating normal to and along the two-dimensional 2D layer. Two extreme intersubband transitions are shown in Fig. 1(a): transition 1 is caused by the phonons with energy close to level splitting Δ_T where 2D momentum transmission is small. The momentum transmission under transition 2 is on the order of $\sqrt{2m\Delta_T}$ and the energy of the phonon involved is on the order of $\sqrt{2ms^2\Delta_T}$ (here m is the electron effective mass and s is the sound velocity). For these transitions we have strong nonmonotonic energy and angle dependencies in the electron-phonon matrix element due to the interwell interference of the electron states of the left (l-) and of the right (r-) QW's.

In this paper we consider linear, with respect to a phonon flux, modulation of the electron response to nonequilibrium phonon flux which could be measured as (i) transverse voltage which arises due to the electron redistribution between QW's, and (ii) phonon drag effect caused by the transfer of in-plane phonon momentum components to the electrons in DQW's. Note that the phonon-induced voltage can be observed for a phonon pump with isotropic distribution of momentum in the 2D plane as shown in Fig. 1(b) while the phonon drag current can only be observed for excitation by the phonon beam with a nonzero total in-plane component of the wave vector as is shown in Fig. 1(c). To the best of our knowledge, the phonon-induced voltage has not been investigated in 2D electron systems. The phonon drag effect could be observed in the geometry of Fig. 1(c) (analogous geometry was realized for photon drag; see Ref. 14 and references therein). The phonon drag effect in selectively doped GaAs heterostructure has been also experimentally investigated.[14,15] Phonon-induced modification of exciton spectra in DQW's was reported last year,[16] whereas experimental study of the ''phonoelectric'' phenomena was not carried out yet. The prime objects of this paper are to describe the main peculiarities of this type of response and to provide the numerical results in addition to our analytical calculations.

In our calculations we use the two-level DQW model that

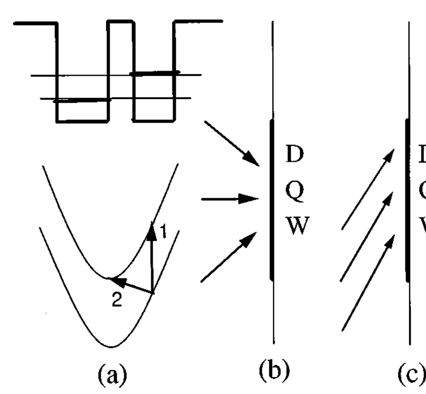

FIG. 1. Energy-band diagram of DQW's and phonon-induced transitions for the cases of normal-incident (1) and ''slipping'' (2) phonon modes (a) and the schemes of excitation under consideration: by the in-plane isotropic pump (b) and by the phonon beam with fixed longitudinal momentum (c).

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was described in Refs. 6 and 17 and its band diagram is shown in Fig. 1(a). The matrix elements for deformation mechanism of phonon interaction with tunnel-coupled electron states in such DQW's were discussed in Ref. 13. We assume that momentum relaxation and the interwell electron redistribution in DQW's are controlled by interface roughness scattering, which, in the general case, can be different for l- and r-QW's. We describe nonequilibrium phonons by means of the phonon population number $\delta N_{\mathbf{q}}$. The isotropic phonon excitation with fixed temperature, which may be estimated from the experimental data, is used for case (i) and the excitation by a phonon beam with fixed wave vector $\mathbf{Q}$ is used for case (ii) (for more details see Ref. 18 where the typical experimental conditions are discussed).

The paper is organized as follows. In Sec. II we present the basic expression for linear response of electrons in DQW's due to phonon perturbation. Phonon-induced voltage (including the special case of identical QW's with nonsymmetric scattering) and phonon-drag effects are considered in Secs. III and IV, respectively. Concluding remarks are given in Sec. V.

II. BASIC EQUATIONS

In this section we consider the linear response of electrons to excitation by ballistic acoustic phonons. Under the description of electron response the calculations below use the electron-phonon collision integral in the quantum kinetic equation. This equation takes into account the interwell coupling of the electron states in the framework of the two-level tunnel approximation.

For a description of the tunnel-coupled ground states in DQW's, we use a basis of the orbitals corresponding to l- and r-QW's, where the Hamiltonian is given by the 2×2 matrix (''isospin'' representation[6])

$$\hat{h}_p=\varepsilon_p+\hat{h}_\perp\,, \quad \hat{h}_\perp=\frac{\Delta}{2}\,\hat{\sigma}_z+T\hat{\sigma}_x\,, \tag{1}$$

where $\varepsilon_p=p^2/2m$ is the kinetic energy of the longitudinal motion ($\mathbf{p}$ is 2D momentum and m is effective mass), Δ is the level splitting, T is the tunnel matrix element, and $\hat{\sigma}_{x,y,z}$ are the Pauli matrixes. The density of the phonon-induced current $\delta\mathbf{J}$ and voltage δU normalized to the unit area L^2 are given by the expressions

$$\delta\mathbf{J}=\frac{e}{m}\frac{2}{L^2}\sum_{\mathbf{p}}\mathbf{p}tr\,\delta\hat{f}_{\mathbf{p}}\,,\quad \delta U=\frac{4\pi e^2Z}{\varepsilon}\frac{2}{L^2}\sum_{\mathbf{p}}tr\,\hat{\sigma}_z\delta\hat{f}_{\mathbf{p}}\,, \tag{2}$$

where tr corresponds to the summation of diagonal matrix elements, Z is the distance between the centers of the orbitals corresponding to l- and r-QW's. Since we are interested in the modulation of the electron response of DQW's electrons with respect to the stationary excitation of acoustic phonons, we will only keep in Eq. (2) the nonequilibrium component of the electron distribution $\delta\hat{f}_{\mathbf{p}}$ caused by this perturbation. Therefore, the 2×2 matrix $\delta\hat{f}_{\mathbf{p}}$ can be determined from the linearized kinetic equation

$$\frac{i}{\hbar}[\hat{h}_\perp,\delta\hat{f}_{\mathbf{p}}]=I(\delta\hat{f}|\mathbf{p})+\delta\hat{I}_{\rm AC}(\mathbf{p}), \tag{3}$$

where $I(\delta\hat{f}|\mathbf{p})$ is the collision integral caused by elastic scattering mechanisms and $\delta\hat{I}_{\rm AC}(\mathbf{p})$ is the perturbation due to acoustic phonon excitation.

Considering the electron collision in Eq. (3) we restrict ourselves to the case of elastic scattering due to interface roughness; such mechanisms have been applied in the description of longitudinal transport and interwell tunneling relaxation.[6] The collision integral, which takes into account nonsymmetrical scattering in l- and r-QW's, is given by

$$\begin{aligned}I(\delta\hat{f}|\mathbf{p})=&\frac{1}{\hbar^2}\sum_{k=l,r}w_k\sum_{\mathbf{p}'}e^{-|\mathbf{p}-\mathbf{p}'|^2(l_k/2\hbar)^2}\\&\times\int_{-\infty}^{0}d\tau\,e^{\lambda\tau}\{\hat{S}_{\mathbf{p}\tau}(\hat{P}_k\delta\hat{f}_{\mathbf{p}'}-\delta\hat{f}_{\mathbf{p}}\hat{P}_k)\\&\times\hat{S}^{+}_{\mathbf{p}'\tau}\hat{P}_k-\hat{P}_k\hat{S}_{\mathbf{p}'\tau}(\hat{P}_k\delta\hat{f}_{\mathbf{p}'}-\delta\hat{f}_{\mathbf{p}}\hat{P}_k)\hat{S}^{+}_{\mathbf{p}\tau}\},\end{aligned} \tag{4}$$

where $\lambda\to+0$, $\hat{S}_{\mathbf{p}\tau}=\exp(i\hat{h}_{\mathbf{p}}\tau/\hbar)$, l_k is the correlation length of the roughness in k-QW's, parameter $w_k=2\pi l_k^2(2\varepsilon_k\delta/d_k)^2$ determines a scattering efficiency in k-QW's, and $\hat{P}_k$ is the projective operator into states of k-QW's (δ is an average roughness amplitude and ε_k is the energy of the ground state in k-QW's with width d_k).

In order to describe the electron response in DQW's it is convenient to use the basis of the tunnel-coupled states that can be determined from the eigenstate problem $\hat{h}_p|\pm\rangle=\varepsilon_{\pm p}|\pm\rangle$ where $\varepsilon_{\pm p}=\varepsilon_p\pm\Delta_T/2$ describes dispersion relations for $|\pm\rangle$ states with splitting energy $\Delta_T=\sqrt{\Delta^2+4T^2}$. According to Eq. (3) the diagonal elements of the electron distribution in this basis $\delta f_{j\mathbf{p}}=\langle j|\delta\hat{f}_{\mathbf{p}}|j\rangle$ are determined by the system of equations for $|\pm\rangle$ states

$$I(\delta f|j\mathbf{p})+\delta I_{\rm AC}(j\mathbf{p})=0. \tag{5}$$

A similar equation may be written for the nondiagonal components $\widetilde{\delta f}_{\mathbf{p}}=\langle+|\delta\hat{f}_{\mathbf{p}}|-\rangle=\langle-|\delta\hat{f}_{\mathbf{p}}|+\rangle^*$, which are small compared to $\hbar/(\Delta_T\tau)$ (τ corresponds to effective scattering time for collision integral). Then we can express $\widetilde{\delta f}_{\mathbf{p}}$ by means of nondiagonal elements of the collision integral (this term accounts for the diagonal contributions of $\delta f_{j\mathbf{p}}$ only) and phonon perturbation $\delta I_{\rm AC}$:

$$\widetilde{\delta f}_{\mathbf{p}}=-\frac{i\hbar}{\Delta_T}[\langle+|I(\delta\hat{f}|\mathbf{p})|-\rangle+\langle+|\delta I_{\rm AC}(\mathbf{p})|-\rangle]. \tag{6}$$

In the basis of $|\pm\rangle$ states the phonon-induced current $\delta\mathbf{J}$ is determined by diagonal elements of the nonequilibrium electron distribution

$$\delta\mathbf{J}=\frac{e}{m}\frac{2}{L^2}\sum_{j\mathbf{p}}\mathbf{p}\delta f_{j\mathbf{p}}. \tag{7}$$

The induced voltage δU takes the form

$$\begin{aligned}\delta U=\frac{4\pi e^2Z}{\varepsilon}\frac{2}{L^2}\sum_{\mathbf{p}}\Bigg[&\frac{\Delta}{\Delta_T}(\delta f_{+\mathbf{p}}-\delta f_{-\mathbf{p}})\\&-\frac{2T}{\Delta_T}(\widetilde{\delta f}_{\mathbf{p}}+\widetilde{\delta f}^{*}_{\mathbf{p}})\Bigg],\end{aligned} \tag{8}$$

and Eq. (8) also contains nondiagonal contributions that become the principal contributions in the tunnel resonance (near $\Delta=0$).

For elastic scattering by interface roughness the diagonal elements of collision integral in Eq. (5) depend both on intersubband and intrasubband electron transition probabilities, $W(j\mathbf{p},j'\mathbf{p}')$, and

$$I(\delta f|j\mathbf{p})=\sum_{j'\mathbf{p}'} W(j\mathbf{p},j'\mathbf{p}')(\delta f_{j'\mathbf{p}'}-\delta f_{j\mathbf{p}}), \tag{9}$$

$$W(j\mathbf{p},j'\mathbf{p}')=\frac{2\pi}{\hbar}\sum_{k=l,r} w_k e^{-|\mathbf{p}-\mathbf{p}'|^2(l_k/2\hbar)^2}\Phi^{(k)}_{j,j'} \times\delta(\varepsilon_{jp}-\varepsilon_{j'p'}).$$

Here the overlap factors are $\Phi^{(l)}_{+,+}=\Phi^{(r)}_{-,-}=(1+\Delta/\Delta_T)^2/4$, $\Phi^{(r)}_{+,+}=\Phi^{(l)}_{-,-}=(1-\Delta/\Delta_T)^2/4$, $\Phi^{(l)}_{+,-}=\Phi^{(r)}_{+,-}=\Phi^{(l)}_{-,+}=\Phi^{(r)}_{-,+}=(T/\Delta_T)^2$. In contrast to Eq. (9), the nondiagonal elements of the collision integral are determined not only by transition probabilities but also by the modification of electron states in DQW's and will be considered in Sec. III B.

In order to complete formulation of the problem to be solved, the acoustic phonon perturbation, $\delta I_{\mathrm{AC}}(\mathbf{p})$ in Eq. (3) [and in Eqs. (5) and (6)], should be expressed through the nonequilibrium phonon distribution $\delta N_{\mathbf{Q}}$ (the equilibrium phonons are insignificant as a source of relaxation because of roughness scattering). The phonon perturbation $\delta\hat{I}_{\mathrm{AC}}(\mathbf{p})$ in such an approximation may be presented as the sum of contributions from emission and absorption of phonons. Equation (10) presents the contribution of the phonon emission into the perturbative term δI_{AC}, and the contribution from phonon-absorption can be obtained from Eq. (10) by substitution $\omega_Q\rightarrow-\omega_Q$ and $\mathbf{q}\rightarrow-\mathbf{q}$,

$$\delta\hat{I}_{\mathrm{AC}}(\mathbf{p})=\frac{1}{\hbar^2}\sum_{\mathbf{Q}}|C_Q|^2\chi(q_\perp d)^2\delta N_{\mathbf{Q}}\int_{-\infty}^{0}d\tau\, e^{\lambda\tau}\{e^{i\omega_Q\tau}[\hat{S}_{\mathbf{p}\tau}(1-\hat{f}_{\mathbf{p}})\hat{\chi}^+_{q_\perp}\hat{f}_{\mathbf{p}+\hbar\mathbf{q}}\hat{S}^+_{\mathbf{p}+\hbar\mathbf{q}\tau}\hat{\chi}_{q_\perp}-\hat{\chi}_{q_\perp}\hat{S}_{\mathbf{p}-\hbar\mathbf{q}\tau}(1-\hat{f}_{\mathbf{p}-\hbar\mathbf{q}})\hat{\chi}^+_{q_\perp}\hat{f}_{\mathbf{p}}\hat{S}^+_{\mathbf{p}\tau}]$$
$$-e^{-i\omega_Q\tau}[\hat{S}_{\mathbf{p}\tau}\hat{f}_{\mathbf{p}}\hat{\chi}_{q_\perp}(1-\hat{f}_{\mathbf{p}-\hbar\mathbf{q}})\hat{S}^+_{\mathbf{p}-\hbar\mathbf{q}\tau}\hat{\chi}^+_{q_\perp}-\hat{\chi}^+_{q_\perp}\hat{S}_{\mathbf{p}+\hbar\mathbf{q}\tau}\hat{f}_{\mathbf{p}+\hbar\mathbf{q}}\hat{\chi}_{q_\perp}(1-\hat{f}_{\mathbf{p}})\hat{S}^+_{\mathbf{p}\tau}]+[\omega_Q\rightarrow-\omega_Q,\mathbf{q}\rightarrow-\mathbf{q}]\}, \tag{10}$$

where C_Q is the bulk matrix element for electron-phonon interaction, $\mathbf{Q}=(\mathbf{q},q_\perp)$ and ω_Q are, respectively, the phonon wave vector and the phonon energy, and $\hat{f}_{\mathbf{p}}$ is the electron-density matrix in the absence of nonequilibrium phonons. The form factors $\chi(q_\perp d)$ and $\hat{\chi}_{q_\perp}$ describing both isolated QW's (the potentials for l- and r-QW's with width $d\approx d_{l,r}$ are assumed to be equal) and the interwell interference effects[13] take the form

$$\hat{\chi}_{q_\perp}=\cos(q_\perp Z)-i\hat{\sigma}_z\sin(q_\perp Z), \quad \chi(a)=\frac{(2/a)\sin(a/2)}{1-(a/2\pi)^2}. \tag{11}$$

Note that the intrawell factor $\chi(a)$ is presented here for the confined QW potential in flat-band approximation (see Ref. 19). For the calculation of the phonon excitation in Eq. (5) we have to rewrite the operator $\delta I_{\mathrm{AC}}(j\mathbf{p})$ in terms of diagonal and nondiagonal (with respect to the above-introduced $\pm$ bases) matrix elements. Taking into account the diagonal elements of the density matrix $\hat{f}$ (which correspond to the distribution functions $f_{\pm,\mathbf{p}}$) and neglecting the small nondiagonal contributions we obtain the general expression for the generation term in Eq. (5),

$$\delta I_{\mathrm{AC}}(j\mathbf{p})=\frac{2\pi}{\hbar}\sum_{j'=\pm}\sum_{\mathbf{Q}}|C_Q|^2\chi(q_\perp d)^2|\langle j|\hat{\chi}_{q_\perp}|j'\rangle|^2\delta N_{\mathbf{Q}} \times(f_{j'\mathbf{p}+\hbar\mathbf{q}}-f_{j\mathbf{p}})[\delta(\varepsilon_{j'\mathbf{p}+\hbar\mathbf{q}}-\varepsilon_{jp}-\hbar\omega_Q) +\delta(\varepsilon_{j'\mathbf{p}+\hbar\mathbf{q}}-\varepsilon_{jp}+\hbar\omega_Q)]. \tag{12}$$

The phonon population number $\delta N_{\mathbf{Q}}$ in Eq. (12) is determined by both the character of the momentum distribution (the in-plane isotropic excitation and the phonon beam with fixed momentum are used in Secs. III and IV) and the energy flux density $\mathbf{G}$. This value is introduced as

$$\mathbf{G}=L^{-3}\sum_{\mathbf{Q}}(\nabla\omega_Q)\hbar\omega_Q\delta N_{\mathbf{Q}}, \tag{13}$$

and Eq. (13) also depends on the spectral distribution of phonons. Below we use the symmetric over 2D-plane phonon distribution and phonon beam with fixed wave vector $\mathbf{Q}$ (in Secs. III and IV, correspondingly).

III. INDUCED VOLTAGE

According to Eq. (8), the phonon-induced voltage over DQW's contains the first term (denoted below as δU_Δ) which is caused by the redistribution of electrons between l- and r-QW's and also the collision-induced modification of the electron states in symmetric DQW's, the second term (denoted below as δU_s) described by the nondiagonal components of distribution function (6). Below we consider these two contributions for the weak-coupled QW's by using a quasiequilibrium distribution of electrons over $\pm$ states.

A. Response of nonsymmetric DQW's

For the low-temperature case, the quasiequilibrium electron distribution $f_{\pm p}$ may be written as $\theta(\mu_\pm-\varepsilon_{\pm p})$, $\mu_\pm$ are the Fermi levels. The expression for δU_Δ may be written as

$$\delta U_\Delta = \frac{4\pi e^2 Z}{\varepsilon} = \frac{\Delta}{\Delta_T}(n_+ - n_-),$$

$$n_\pm = \frac{2}{L^2}\sum_{\mathbf{p}} \theta(\mu_\pm - \varepsilon_{\pm p}), \tag{14}$$

and the induced voltage is proportional to the redistribution of the electron concentration between ''+'' and ''−'' states $\delta n = n_+ - n_-$. After summing the system of Eqs. (5) over $\mathbf{p}$, we obtain the balance equations for the concentrations $n_\pm$. For the case of a small redistribution of the concentration [if $\delta n = \ll (n_+ + n_-)/2$], the equation for the phonon-induced redistribution transforms to the form

$$\delta I^+_{\rm AC} + \nu_T \delta n = 0. \tag{15}$$

The rate of interwell tunneling ν_T is expressed through the probability of transitions (9) according to the relation

$$\nu_T(p) = \sum_{\mathbf{p}'} W(+\mathbf{p}, -\mathbf{p}'), \tag{16}$$

and the effective tunneling rate in Eq. (15), ν_T, is presented and discussed in Ref. 20 for the scattering due to interface roughness. The concentration dependence of ν_T is weak but the dependence of ν_T on the level splitting Δ is substantial and below we assume the dependence $\nu_T = \overline{\nu_T}(2T/\Delta_T)^2$; $\overline{\nu_T}$ is the resonant tunneling rate.

The phonon-induced contribution to Eq. (15), $\delta I_{\rm AC}$, does not depend on δn for the small redistribution case under consideration and this contribution may be transformed from Eqs. (11) and (12) to the form

$$\begin{aligned}\delta I^+_{\rm AC} = &\frac{(2T/\Delta_T)^2}{\tau_{\rm ph} T}\int_0^\infty d\varepsilon \int_0^\infty d\varepsilon' \int_0^{2\pi}\frac{d\phi}{2\pi}\int dq_\perp \chi(q_\perp d)^2 \\ &\times \sin^2(q_\perp Z)\,\delta N_Q Q\left[\theta\left(\mu + \frac{\Delta_T}{2} - \varepsilon\right)\right. \\ &\left. - \theta\left(\mu - \frac{\Delta_T}{2} + \varepsilon\right)\right][\delta(\varepsilon' - \varepsilon - \Delta_T - \hbar c Q) \\ &+ \delta(\varepsilon' - \varepsilon + \Delta_T - \hbar c Q)].\end{aligned} \tag{17}$$

Here we have introduced the phonon wave number Q and the characteristic relaxation time $\tau_{\rm ph}$ according to the relations

$$\hbar Q = \sqrt{2m[\varepsilon + \varepsilon' - \sqrt{\varepsilon\varepsilon'}\,\cos\phi] + (\hbar q_\perp)^2},$$

$$\frac{1}{\tau_{\rm ph}} = \frac{2D^2 m^2 T}{\pi \hbar^4 \rho s}, \tag{18}$$

where D is the deformation potential constant, ρ is heterostructure density, and the angle $\phi = (\mathbf{p}, \hat{\mathbf{p}}')$ appears due to the in-plane symmetry of the problem. For the case of high-energy excitation of DQW's we use below the phonon distribution $\delta N_Q = [\exp(\hbar\omega_Q/T_{\rm ph}) - 1]^{-1}$, for $q_\perp > 0$ and $\delta N_Q = 0$ for $q_\perp < 0$; $T_{\rm ph}$ is the phonon temperature. Integration over ϕ may be performed by using δ functions and the integral over $q_\perp$ may be taken analytically under the condition $\sqrt{2ms^2\Delta_T} \ll T_{\rm ph}$, when the narrow region of the transverse wave vectors gives main contribution to the integrand. The last double integral (over ε and ε') may be evaluated numerically and the results for δU_Δ are obtained from Eqs. (14) and (15).

In order to discuss the numerical results for δU_Δ, we consider a typical symmetrical GaAs/$\mathrm{Al_{0.3}Al_{0.7}}$As-based DQW structure with the following parameters: $d \approx 70$ Å, $Z \approx 120$ Å; the barrier width is equal to 47, 56, and 65 Å (these values correspond to the tunnel matrix elements 2, 1, and 0.5 meV), the tunneling relaxation rate for the near-resonant region $\overline{\nu_T}$ is estimated below as 50 ps and the level-splitting dependency of $\nu_T(\Delta)$ is taken into account (see above). We use $T_{\rm ph} = 2$ meV for the high-energy phonon pump here. In Figs. 2(a)–2(c) we plot the dependencies of the phonon-induced voltage on the level splitting Δ, which can be controlled by the application of a transverse gate voltage (so that we can consider Δ as an independent parameter). The voltage ΔU_Δ vanishes for both small and large Δ (due to symmetry of DQW's if $\Delta = 0$ and due to suppression of tunnel coupling for $\Delta \gg 2T$, correspondingly) and the amplitude of δU_Δ peak increases with concentration. The voltage decreases if T increases. The shift in voltage maximum with increased T is demonstrated in Fig. 3 for the normalized δU_Δ. If $T_{\rm ph}$ decreases the peak of δU_Δ is shifted by a small Δ_T and the sensitivity of the phonon-induced response is limited by the scattering processes or by the large-scale inhomogeneities of the DQW structure.

B. Voltage due to nonsymmetric scattering

Below, we consider the case of symmetric DQW's ($\Delta = 0$) since the small contribution due to nonsymmetrical scattering is essential for only small Δ. The nondiagonal contributions, both from the collision integral (4) and from the acoustic-phonon perturbation (10) should be considered in Eq. (6). Using the diagonal electron contribution $f_{\pm p}$ in Eq. (10) and the matrix elements of the operator (11) we have found that the matrix element in $\langle + | \delta I_{\rm AC}(\mathbf{p}) | - \rangle$ is proportional to $\cos q_\perp Z \sin q_\perp Z$. Since $\chi(a) = \chi(-a)$, then the phonon contribution to $\widetilde{\delta f}_{\mathbf{p}}$ is equal to zero. As a result, the nondiagonal component of $\widetilde{\delta f}_{\mathbf{p}}$ is written in the form

$$\begin{aligned}\widetilde{\delta f}_{\mathbf{p}} = &\frac{\mathcal{P}}{i\Delta_T}\int \frac{d\mathbf{p}_1}{(2\pi\hbar)^2}\sum_k w_k \exp[-|\mathbf{p} - \mathbf{p}'|^2(l_k/2\hbar)^2] \\ &\times \sum_j \langle + | \hat{P}_k | j \rangle \langle j | \hat{P}_k | - \rangle \left[\frac{\delta f_{-\mathbf{p}} - \delta f_{j\mathbf{p}_1}}{\varepsilon_{jp} - \varepsilon_{-p}}\right. \\ &\left. - \frac{\delta f_{+\mathbf{p}} - \delta f_{j\mathbf{p}}}{\varepsilon_{jp} - \varepsilon_{+p}}\right],\end{aligned} \tag{19}$$

where $\mathcal{P}$ means the principal value. Using the definition (8) we get the following phonon-induced voltage:

$$\begin{aligned}\delta U_s = &\frac{8\pi e^2 Z}{\varepsilon T}\mathcal{P}\int \frac{d\mathbf{p}}{(2\pi\hbar)^2}\int \frac{d\mathbf{p}_1}{(2\pi\hbar)^2}\frac{\delta f_{-p} - \delta f_{+p}}{\varepsilon_p - \varepsilon_{p_1}} \\ &\times \{w_l \exp[-|\mathbf{p} - \mathbf{p}'|^2(l_l/2\hbar)^2] - w_r \\ &\times \exp[-|\mathbf{p} - \mathbf{p}'|^2(l_r/2\hbar)^2]\}.\end{aligned} \tag{20}$$

Note that for the symmetric scattering case ($w_l = w_r$ and $l_l = l_r$) this contribution vanishes.

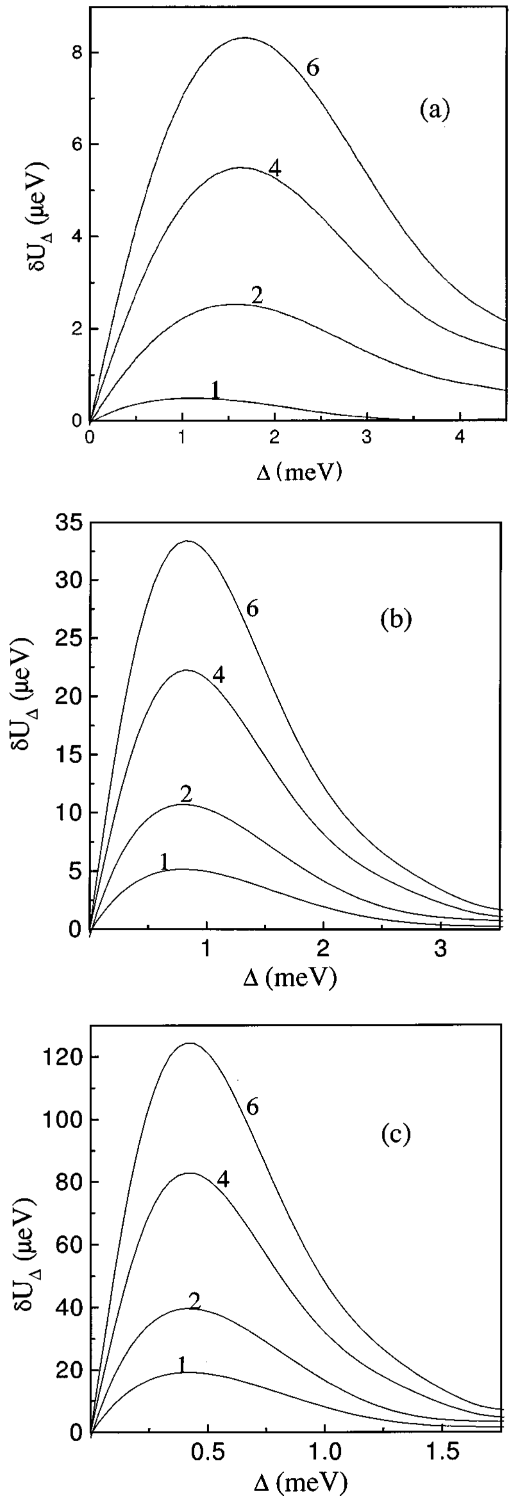

FIG. 2. The dependence of the phonon-induced voltage δU_Δ vs the level splitting Δ for DQW's with the tunnel matrix elements $T=2$ meV (a), 1 meV (b), and 0.5 meV (c). The curves marked by the indexes 1, 2, 4, and 6 correspond to the electron concentrations 10^{11}, 2×10^{11}, 4×10^{11}, and 6×10^{11} cm^{-2}, respectively.

A simple estimation for Eq. (20) results for the case of short-range scattering ($l_k\rightarrow0$), when the integration over momentum is restricted by the condition $\varepsilon_p,\varepsilon_{p_1}<\varepsilon_m$ [the cutoff energy ε_m is estimated as $(\hbar/l_k)^2/2m$]. The logarithmically divergent (if $l_k\rightarrow0$) contributions appear after the integration in Eq. (20),

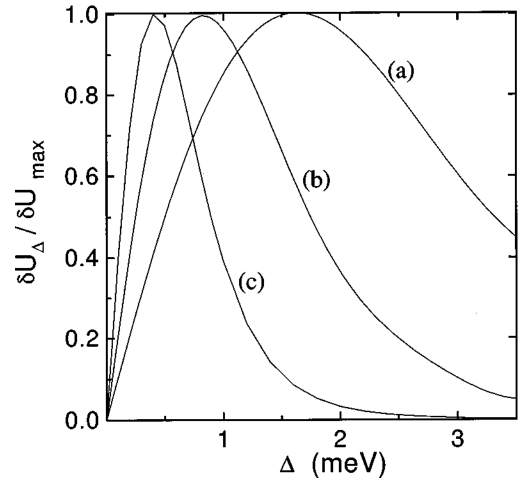

FIG. 3. Normalized level-splitting dependencies of the induced voltage for the DQW's with different tunnel matrix elements: $T=2$ meV (a), 1 meV (b), and 0.5 meV (c).

$$\mathcal{P}\int\frac{d\mathbf{p}}{(2\pi\hbar)^2}(\varepsilon_p-\varepsilon)^{-1}=\frac{\rho_{2\mathrm{D}}}{2}\ln\frac{\varepsilon_m}{\varepsilon},\tag{21}$$

and analogous with the single QW case[21] we can introduce a large logarithmical factor $\Lambda=\ln(\varepsilon_m/\overline{\varepsilon})$; $\overline{\varepsilon}$ is characteristic electron energy. The second integral gives the redistribution of concentration δn and the result takes the form

$$\delta U_s=\frac{4\pi e^2Z}{\epsilon}\delta n\Lambda\frac{\hbar}{2T}(\tau_l^{-1}-\tau_r^{-1}),\tag{22}$$

where $\tau_k=\pi\rho_{2\mathrm{D}}w_k/\hbar$ is the relaxation time in the kth QW. In comparison with the contribution from δU_Δ (which is proportional to Δ/Δ_T) the quantum factor $\Lambda(\hbar/2T)(\tau_l^{-1}-\tau_r^{-1})$ appears in Eq. (22). The contribution of Eq. (22) is larger than δU_Δ if

$$\Delta<\Lambda\hbar(\tau_l^{-1}-\tau_r^{-1}),\tag{23}$$

i.e., for the narrow near-resonant region.

For the estimations of the characteristic splitting levels when Eq. (23) is valid, we use $\Lambda\approx5$ and take the mobilities in l- and r-QW's equal to 2×10^5 and 10^6 cm^2/V s, respectively. As a result, this contribution is essential when Δ is smaller then 0.2 meV for the DQW's with the above-listed parameters and δU_s is on the order of 1 μeV or smaller.

IV. PHONON DRAG CURRENT

The momentum transfer under phonon excitation of DQW's is responsible for the drag current (7), which is expressed through the asymmetrical part of the electron distribution function, $\delta f_{j\mathbf{p}}^{\mathrm{as}}$. Equations (5) and (6) give us

$$\delta f_{j\mathbf{p}}^{\mathrm{as}}=\tau_{jp}\delta I_{\mathrm{AC}}^{\mathrm{as}}(j\mathbf{p}),\tag{24}$$

where τ_{jp} is the momentum relaxation time and $\delta I_{\mathrm{AC}}^{\mathrm{as}}(j\mathbf{p})$ is the asymmetric part of Eq. (12). Below we consider short-range roughness scattering and restrict ourselves to the case of $\mu>\Delta_T/2$ (when $\pm$ electron states are filled). Under these conditions the momentum relaxation rate takes the form

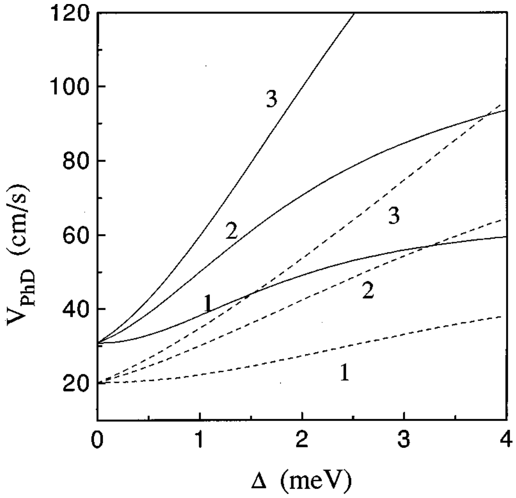

FIG. 4. The level-splitting dependencies of the phonon-drag velocity $V_{\rm PhD}$ for the DQW's with the tunnel matrix elements $T=2$ meV (dashed) and $T=1$ meV (solid). The factors of scattering asymmetry: $\tau_l/\tau_r=1$ (1), $\tau_l/\tau_r=0.3$ (2), and $\tau_l/\tau_r=0.1$ (3).

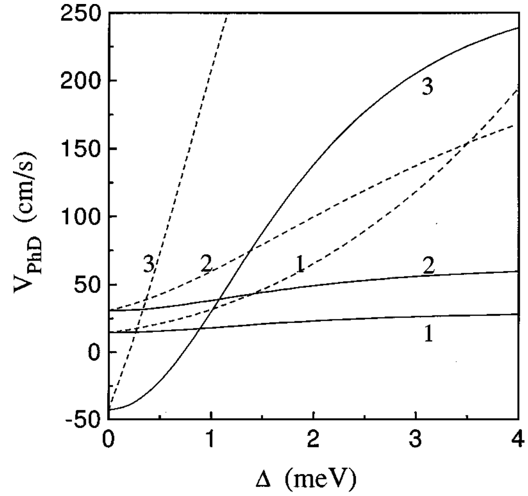

FIG. 5. The same as in Fig. 4 for the scattering asymmetry factor $\tau_l/\tau_r=1$ (solid) and $\tau_l/\tau_r=0.1$ (dashed) for the incidence angles: $\theta=60°$ (1), $\theta=45°$ (2), and $\theta=30°$ (3).

$$\tau_{\pm}^{-1}=\frac{\tau_l^{-1}+\tau_r^{-1}}{2}\pm\frac{\Delta}{\Delta_T}\,\frac{\tau_l^{-1}-\tau_r^{-1}}{2}, \tag{25}$$

where the second term stands for the scattering asymmetry. For the case of excitation, the phonon beam $\delta I_{\rm AC}^{\rm as}(j\mathbf{p})$ transforms to

$$\begin{aligned}\delta I_{\rm AC}^{\rm as}(j\mathbf{p})=&\frac{\Omega_0}{4\hbar}\sum_{j'=\pm}\sum_{\mathbf{Q}}|C_Q|^2\chi(\widetilde{q}_\perp\overline{d})^2|\langle j|\hat{\chi}_{\widetilde{q}_\perp}|j'\rangle|^2\delta N_Q\\&\times(f_{j'\mathbf{p}+\hbar\mathbf{q}_\perp}-f_{j\mathbf{p}})[\delta(\varepsilon_{j'\mathbf{p}+\hbar\mathbf{q}_\parallel}-\varepsilon_{jp}-\hbar\omega_Q)\\&+\delta(\varepsilon_{j'\mathbf{p}+\hbar\mathbf{q}_\parallel}-\varepsilon_{jp}+\hbar\omega_Q)]+(\mathbf{p}\rightarrow-\mathbf{p}),\end{aligned} \tag{26}$$

where Ω_0 is the angular width of the beam propagating along the vector with fixed direction $(\mathbf{q}_\parallel,\widetilde{q}_\perp)$. For the incidence angle θ we use $q_\parallel=Q\sin\theta$, $\widetilde{q}_\perp=Q\cos\theta$ and one-dimensional integral over Q stays in Eq. (26) only. It should be noted that $\delta I_{\rm AC}^{\rm as}$ is not a positive definite value. Due to the phonon-induced intersubband transitions the inversion of $V_{\rm PhD}$ for small Δ (which is similar to the spectral inversion of photon drag effect; see Ref. 14) is possible.

Now, we introduce the phonon-induced drift velocity $V_{\rm PhD}$ according to the relation

$$j=enV_{\rm PhD}, \tag{27}$$

where n is the total 2D concentration of electrons. Using Eq. (7) and the above-presented relations, we obtain the numerical results for $V_{\rm PhD}$. To be specific, we will consider the DQW's with parameters introduced in Sec. III and with 2D concentration $n=4\times10^{11}$ cm^{-2}. Figure 4 shows the dependence of the velocity $V_{\rm PhD}$ on level splitting for the different tunnel matrix elements and the different scattering asymmetry factor; the incidence angle here is equal to 45°. The functions $V_{\rm PhD}(\Delta)$ for the different incidence angles and tunnel matrix elements are presented in Fig. 5. The most interesting property of these dependencies is the origin of the *inversion* of the phonon-drag current. For small Δ (when tunnel mixing is essential) and for short-range scattering regimes, the interlevel phonon-induced transitions may be dominant. Under these conditions the momentum transfer for $\pm$ states is near $\pm\mathbf{q}/2$ and the total current may change direction. Such behavior is analogous to the spectral inversion of interband photon drag current.[14]

V. CONCLUSION

In this paper, we have considered two types of the phonon-induced electron responses in DQW's (transverse voltage and phonon-drag current) and have presented the dependencies of δU and $V_{\rm PhD}$ versus level splitting, concentration, and nonsymmetry of scattering. These results demonstrate large values for the phenomena under consideration. For typical DQW parameters, the characteristic phonon-induced potential is on the order of 10 μeV under high-energy phonon pump intensity on the order of 1 μW/cm^2. For the DQW's symmetrical level-induced voltages (typical value is on the order of 1 μV) appear due to scattering asymmetry. The phonon-induced drag effect is characterized by the velocity $V_{\rm PhD}$ which is on the order of 100 cm/s under the excitation by a beam with an intensity of approximately 1 nW and an angle width of 10°. The sign of the drag velocity may be changed in the vicinity of the tunneling resonance where interlevel transitions are essential.

Our theory was developed under the following assumptions. The model of electron states use a simple two-level approach and does not take into account self-consistent corrections, so Δ and T in the matrix Hamiltonian (1) are the semiphenomenological parameters. Using the Planck function as a nonequilibrium phonon-energy distribution we may consider the effective phonon temperature $T_{\rm ph}$ as a characteristic of pump intensity only. The model of the scattering processes due to interface roughness describes the principal properties of both momentum and interwell relaxation. Consideration of the linear response case is justified as long as δU and $V_{\rm PhD}$ are small in comparison with the level splitting Δ_T and the corresponding Fermi velocity. The further restriction is due to the deformation electron-phonon interaction only; this contribution is dominant for high-energy excitation in comparison with piezoelectric interaction. The above assumptions do not change significantly the presented estimations for δU and V_{PhD} or their dependencies on the DQW parameters.

The above-discussed phenomena are sensitive enough on the parameters of phonon excitation (i.e., energy and angle distributions of phonon) as long as δU and V_{PhD} strongly depend on level splitting Δ and incident angle θ. The phonon-modulation of the longitudinal conductivity in DQW's with asymmetric scattering is also well pronounced (such methods have been used for other systems; see e.g., Refs. 22–24) due to the substantial temperature dependence of the conductivity for DQW's with asymmetric scattering. Therefore, presented estimations of phonon-induced responses demonstrate the possibility of phonon detectors based on DQW heterostructures.

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[1] J. P. Eisenstein, Superlattices Microstruct. **12**, 107 (1992).
[2] A. P. Heberle, X. Q. Zhou, and A. Tackeuchi, Semicond. Sci. Technol. **9**, 519 (1994).
[3] A. Palevski, F. Beltram, F. Capasso, L. N. Pfeiffer, and K. W. West, Phys. Rev. Lett. **65**, 1929 (1990).
[4] A. Kurobe, I. M. Castleton, E. H. Linfield, M. P. Grimshaw, K. M. Brown, D. A. Ritchie, M. Pepper, and G. A. C. Jones, Phys. Rev. B **50**, 8024 (1994); **50**, 4889 (1994).
[5] S. G. Boebinger, A. Passner, L. N. Pfeiffer, and K. W. West, Phys. Rev. B **43**, 12 673 (1991).
[6] O. E. Raichev and F. T. Vasko, Phys. Rev. B **53**, 1522 (1996).
[7] Y. S. Huang, H. Qiang, F. H. Pollak, J. Lee, and B. Elman, J. Appl. Phys. **70**, 3808 (1991).
[8] V. Berger, E. Rosencher, N. Vodjdany, and E. Costard, Appl. Phys. Lett. **62**, 378 (1993).
[9] J. Faist, F. Capasso, A. L. Hutchinson, L. Pfeiffer, and K. W. West, Phys. Rev. Lett. **71**, 3573 (1993).
[10] E. Dupont, D. Delacourt, and M. Papuchon, Appl. Phys. Lett. **63**, 2515 (1993).
[11] A. Lorke, A. D. Wieck, U. Merkt, G. Weimann, and W. Schlapp, Semicond. Sci. Technol. **5**, 279 (1989).
[12] J. N. Heyman, K. Unterrainer, K. Craig, B. Galdrikian, M. S. Sherwin, K. Campman, P. F. Hopkins, and A. C. Gossard, Phys. Rev. Lett. **74**, 2682 (1995).
[13] F. T. Vasko, O. G. Balev, and P. Vasilopoulos, Phys. Rev. B **47**, 16 433 (1993).
[14] V. M. Shalaev, C. Douketis, J. T. Stuckless, and M. Moskovits, Phys. Rev. B **53**, 11 388 (1996); F. T. Vasko, *ibid.* **53**, 9576 (1996).
[15] H. Karl, W. Dietsche, A. Fischer, and K. Ploog, Phys. Rev. Lett. **61**, 2360 (1988).
[16] A. V. Akimov, E. S. Moskalenko, L. J. Challis, and A. A. Kaplyanskii, Physica B **219&220**, 9 (1996).
[17] F. T. Vasko and O. E. Raichev, Zh. Eksp. Teor. Fiz. **108**, 2103 (1995) [JETP **81**, 1146 (1995)].
[18] L. J. Challis, in *The Physics of Low-Dimensional Semiconductor Structures*, edited by Butcher, N. H. March, and M. P. Tosi (Plenum, New York, 1993).
[19] P. J. Price, Ann. Phys. (N.Y.) **133**, 217 (1981).
[20] F. T. Vasko and O. E. Raichev, Phys. Rev. B **50**, 12 195 (1994).
[21] C. B. Duke, Phys. Rev. **168**, 816 (1968).
[22] B. Danilchenko, A. Klimashov, O. Sarbey, M. Ashe, and R. Hey, in *23rd International Conference on The Semiconductor Physics*, edited by M. Scheffler and R. Zimmermann (World Scientific, Berlin, 1996), Vol. 3, p. 2339.
[23] J. K. Wigmore, M. Erol, M. Sahraoui-Tahar, C. D. Wilkinson, J. H. Davies, and C. Stanley, Semicond. Sci. Technol. **6**, 837 (1991).
[24] Y. W. Kershaw, S. J. Bending, W. Dietsche, and K. Eberl, Semicond. Sci. Technol. **11**, 1036 (1996).