Control of electron current by double-barrier structures using pulsed laser fields

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Tunneling current through a double-barrier structure is considered in a strong pulsed ac electric field. By absorbing and emitting photons, an electron is able to penetrate through the barriers. An analytical expression for the current is found in the resonance approximation for electric fields with slowly varying amplitudes. The time-dependent evolution of the current exhibits positive and negative amplitudes, depending on the intensity of the field. With an increase of temperature, the current increases substantially. At particular field intensities, a sequence of input laser pulses can be mapped to no pulses, a single pulse, or two pulses in the output electric current. Thus, the decoding of information is essentially impossible.

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I. INTRODUCTION

A remarkable feature of the solid state is the possibility of designing and fabricating artificial structures known as double barriers. These semiconductor structures are formed by alternating ultrathin layers <100 Å of two semiconductor materials that differ in the size of their band gaps, such as in GaAs and $Ga_xAl_{1-x}As$.

These structures are of special interest for several reasons. First, the design of these objects, along with the determination of the appropriate resonance parameters, provides control of electron transport for specific studies and applications. Also, a very interesting feature, very similar to a differential negative resistance, has been found in these structures.[1]

Extensive investigation has been conducted on these double-barrier systems. In 1973 Tsu and Esaki[2] considered a simple model of electron tunneling in such systems, and their theory was later verified experimentally by Chang, Esaki, and Tsu.[3] They found a peak in the current-voltage characteristic corresponding to the electronic level inside the well. By fabricating the DBS with the technique of molecular beam epitaxy, Solner *et al.*[4] provided evidence of resonant tunneling features in current-vs-voltage curves at room temperatures. At 25 K, the peak-to-valley ratio was 6:1. Through improvements in the quality of experimental methods, the structure of $GaAs/Ga_xAl_{1-x}As$ obtained in Ref. 5 had the peak-to-valley ratio 20:1. In certain structures grown by molecular beam epitaxy involving InAlAs/InGaAs,[6] a peak-to-valley ratio as large as 30:1(Ref. 7) has been produced.

When a time-dependent radiation field is applied, strong resonance effects can take place. By absorbing (or emitting) one or two photons, an electron can reach the resonance level and transmit through the DBS. As shown in Ref. 8, the current can have either a positive or negative value, depending upon the intensity of the applied field. The negative amplitude in the current demonstrates the effect of *absolute negative resistance*. Experimentally, the effect of a THz field on electron tunneling was discussed in Ref. 4. Later, absolute negative resistance was found in experiments with a free electron laser.[9]

In a strong laser field, the absorption (or emission) probability of a photon is not small, and the character of tunneling can be changed dramatically. As shown numerically by Hänggi and co-workers,[10] the electron can be localized in the initial well of a double-well system. Also, the unusual behavior of the emission spectra was reported by Bavli and Metiu.[11] They found that, in a symmetric double-well potential, even harmonics are present, while odd harmonics are suppressed. All of these different effects found their explanation in the two electronic level description.[12] By making a proper choice of the laser parameters, it is possible to eliminate selected lines from the spectrum.

When a driven electron interacts with phonons in molecular double-well structures (e.g., electron-transfer chemical reactions), however, the tunneling behavior is qualitatively different. Coherent oscillations in the electron density can be changed to exponential evolution for the electron density. This means that the decay lifetime is dramatically dependent on the laser parameters.[13] All of these effects are indicitave of the unusual properties of *I*-*V* characteristics of quantum structures irradiated by a high intensity light source.[14–19]

In this paper, we study the current in an electric circuit when a double-barrier structure (DBS) is irradiated by a pulsed field. More specifically, we are interested in how the field modulates the absolute positive or negative resistance, and how the current-voltage characteristics depend on laser intensity, temperature, and pulse shape. By varying the pulsed field, we want to investigate the use of an input sequence of optical pulses to encode some desired information, and the employment of an electronic device to transform this information into a sequence of pulses in the electrical circuit.

We, therefore, would like to answer the following basic questions: can the output sequence of current pulses be made to adequately correspond to the input optical information, and is it always possible to decode the output information? To answer these questions, we will use the approach which was employed by Evans *et al.* in Ref. 20 to describe the modulation of the rate constant by an applied pulsed field for the long distance electron transfer rate in polar solvents.

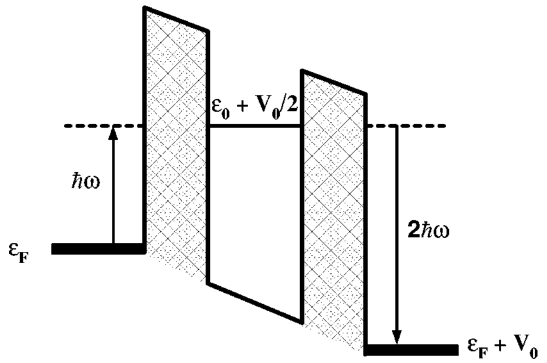

FIG. 1. Double-barrier structure. $\epsilon_F = 0.005$ eV, $n = 10^{17}$ cm^{-3} (Ref. 1). ϵ_0 is the energy of the localized level, ϵ_F is assumed to be much smaller than $\hbar\omega$.

II. RESONANT TRANSMISSION THROUGH DOUBLE-BARRIER STRUCTURE

To find the tunneling current in DBS in an arbitrary time-dependent field, one considers the following Hamiltonian:

$$H = \sum_p \epsilon_p \Psi_p^+ \Psi_p + \sum_k (\epsilon_k + V_0)\Psi_k^+ \Psi_k + (\epsilon_0 + V_0/2 + V(t)/2)\Psi_0^+ \Psi_0 + \sum_p (T_{1p}\Psi_p^+ \Psi_0 + T_{1p}^* \Psi_0^+ \Psi_p) + \sum_k (T_{2k}\Psi_k^+ \Psi_0 + T_{2k}^* \Psi_0^+ \Psi_k). \quad (1)$$

Here, the first and the second terms represent metal leads consisting of highly doped layers of conducting GaAs. The third term describes the position of the local level ϵ_0, inside the well, where V_0 is an applied voltage. Periodically introduced layers of $Ga_x Al_{1-x} As$ serve as barriers with heights of approximately 0.3–0.5 eV. The barrier width ranges from a few tenths to one hundred Å, and to simplify calculations, the barriers are considered to be symmetric. As shown in Fig. 1, ϵ_F, the Fermi energy, is $\epsilon_F = 0.005$ eV for $n = 10^{18}$ cm^{-3},[4] while doping in the center of the well is lower, $n = 10^{17}$ cm^{-3}. At such small concentrations of the electrons, the conduction zones in leads are very narrow, $\epsilon_F \ll \hbar\omega$, where ω is the frequency of the the oscillating time-dependent field [see Eq. (11) below]. According to Solner *et al.*,[4] the position of the electronic level ϵ_0 is about 0.23 eV. The fourth and fifth terms in Eq. (1) represent tunneling through the leads with the transition matrix elements T_{1p} and T_{2k} for the left and the right contacts, respectively. Ψ_0 (Ψ_0^+) is the annihilation (creation) operator inside the well.

A similar Hamiltonian was considered by Ratner and co-workers[21] (with no field) and Tikhonov *et al.*[22] including the time-dependent field, who have applied it to photon-assisted electron tunneling in molecular wires. In Eq. (1) the driving force $V(t)$ is defined as follows:

$$V(t) = \mu_0 E(t), \quad (2)$$

where $E(t)$ is an arbitrary time-dependent electric field and μ_0 is the dipole moment proportional to the width of the DBS. Note that electron-phonon effects are not considered in Hamiltonian (1). Such effects are expected to be small in the temperature regime of primary interest here ($kT \ll \hbar\omega_D$ or $\hbar\omega$, where ω_D is a characteristic phonon frequency and ω is the frequency of the electromagnetic field, respectively).

We seek a solution of the time-dependent Schrödinger equation with Hamiltonian (1). The wave function can be presented in the following form:

$$\Psi(t) = \sum_p \Psi_p(t)|p\rangle + \Psi_0(t)|0\rangle + \sum_k \Psi_k(t)|k\rangle. \quad (3)$$

The Schrödinger equation is then given by the following set of kinetic equations:

$$i\hbar \frac{d\Psi_p}{dt} = \epsilon_p \Psi_p + T_{1p}\Psi_0,$$

$$i\hbar \frac{d\Psi_0}{dt} = [\epsilon_0 + V_0/2 + V(t)/2]\Psi_0 + \sum_p T_{1p}^* \Psi_p + \sum_k T_{2k}^* \Psi_k,$$

$$i\hbar \frac{d\Psi_k}{dt} = (\epsilon_k + V_0)\Psi_k + T_{2k}\Psi_0. \quad (4)$$

Further, we make the following substitution:[8]

$$\Psi_p = \exp\left(-\frac{i}{\hbar}\epsilon_p t\right)\phi_p,$$

$$\Psi_0 = \exp\left[-\frac{i}{\hbar}[(\epsilon_p t + V_0/2)t + F(t)/2]\right]\phi_0,$$

$$\Psi_k = \exp\left[-\frac{i}{\hbar}(\epsilon_k + V_0)t\right]\phi_k. \quad (5)$$

Here, $F(t)$ stands for the following integral:

$$F(t) \equiv \int_0^t dt_1 V(t_1). \quad (6)$$

Thus, we can transform Schrödinger equation (4) to the following form:

$$i\hbar \frac{d\phi_p}{dt} = T_{1p}\exp\left[\frac{i}{\hbar}[(\epsilon_p t - \epsilon_0 t - V_0/2)t - F(t)/2]\right]\phi_0,$$

$$i\hbar \frac{d\phi_0}{dt} = \sum_p T_{1p}^* \exp\left[-\frac{i}{\hbar}[(\epsilon_p t - \epsilon_0 t - V_0/2)t - F(t)f]\right]\phi_p + \sum_k T_{2k}^* \exp\left[-\frac{i}{\hbar}[(\epsilon_k t - \epsilon_0 t + V_0/2)t - F(t)/2]\right]\phi_k,$$

$$i\hbar\frac{d\phi_k}{dt}=T_{2k}\exp\left[\frac{i}{\hbar}[(\epsilon_k t-\epsilon_0 t+V_0/2)t-F(t)/2]\right]\phi_0, \tag{7}$$

with the initial conditions

$$\phi_p(t=0)=1,$$
$$\phi_0(t=0)=0, \tag{8}$$
$$\phi_k(t=0)=0.$$

From the first and the third equations of set (7), one finds that

$$\phi_p=1-iT_{1p}\int_0^t dt_1\phi_0(t_1)\times\exp\left[\frac{i}{\hbar}[(\epsilon_p t-\epsilon_0 t-V_0/2)t-F(t)/2]\right], \tag{9}$$

$$\phi_k=iT_{2k}\int_0^t dt_1\phi_0(t_1)\exp\left[\frac{i}{\hbar}[(\epsilon_k t-\epsilon_0 t+V_0/2)t-F(t)/2]\right].$$

Substituting ϕ_p and ϕ_k from Eqs. (9) into the second equation in Eqs. (7), one arrives at the following integrodifferential equation for $\phi_0(t)$:

$$\hbar\frac{d\phi_0}{dt}=i\sum_p T_{1p}\exp\left[\frac{i}{\hbar}[(\epsilon_0-\epsilon_p+V_0/2)t+F(t)/2]\right]-\int_0^t dt_1\phi_0(t_1)\left\{\sum_p|T_{1p}|^2\exp\left[\frac{i}{\hbar}\{(\epsilon_0-\epsilon_p+V_0/2)\times(t-t_1)+[F(t)-F(t_1)]/2\}\right]+\sum_k|T_{2k}|^2\exp\left[\frac{i}{\hbar}\{(\epsilon_0-\epsilon_k+V_0/2)(t-t_1)+[F(t)-F(t_1)]/2\}\right]\right\}. \tag{10}$$

Furthermore, we consider the slow-varying amplitude of the electromagnetic field $E(t)$ in the cw approximation

$$V(t)=\mu_0 E(t)\cos(\omega t). \tag{11}$$

Thus, $F(t)$ becomes

$$F(t)=\int_0^t dt_1\mu_0 E(t_1)\cos(\omega t_1)\simeq 2a(t)\sin(\omega t), \tag{12}$$

where

$$a(t)\equiv\frac{\mu_0 E(t)}{2\hbar\omega}. \tag{13}$$

Substituting Eq. (13) into Eq. (10) and making use of the identity[23]

$$\exp(\pm ia\sin\omega t)=\sum_{n=-\infty}^{\infty}J_n(a)\exp(\pm in\omega t) \tag{14}$$

we arrive at the following equation for ϕ_0:

$$\hbar\frac{d\phi_0}{dt}=i\sum_p\sum_{m=-\infty}^{\infty}J_m[a(t)]T_{1p}\times\exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2+m\hbar\omega)t\right]-\sum_{m_1=-\infty}^{\infty}\sum_{m_2=-\infty}^{\infty}J_{m_1}[a(t)]\int_0^t dt_1\phi_0(t_1)J_{m_2}[a(t_1)]\times\left\{\sum_p|T_{1p}|^2\exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2)(t-t_1)+im_1\omega t_1-im_2\omega t_2\right]+\sum_k|T_{2k}|^2\exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_k-V_0/2)(t-t_1)+im_1\omega t_1-im_2\omega t_2\right]\right\}. \tag{15}$$

Assuming the tunneling matrix elements T_{1p} and T_{2k} to be small, one notices that the main contribution to the second term in Eq. (15) is due to the resonance or near resonance terms,[8] i.e.,

$$|\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega|\ll\hbar\omega,$$
$$|\epsilon_0-\epsilon_k-V_0/2+n_0\hbar\omega|\ll\hbar\omega, \tag{16}$$

where m_0 and n_0 are determined by resonance conditions (16). In the resonance approximation, Eq. (15) can, therefore, be transformed as follows:

$$\hbar\frac{d\phi_0}{dt}=iJ_{m_0}[a(t)]T_{1p}\exp\left[\left(\frac{i}{\hbar}\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega\right)t\right]-J_{m_0}[a(t)]\int_0^t dt_1\phi_0(t_1)J_{m_0}[a(t_1)]\times\left\{\sum_p|T_{1p}|^2\exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega+i\hbar\delta)\times(t-t_1)\right]+\sum_k|T_{2k}|^2\exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_k-V_0/2+n_0\hbar\omega+i\hbar\delta)(t-t_1)\right]\right\}. \tag{17}$$

In Eq. (17) we have introduced a infinitesimally small imaginary part δ, which determines the elastic scattering of the electron in a DBS. If the system is close to resonance, one can evaluate the integral in Eq. (17) by making use of inverse Laplace transform in the following manner:

$$\int_0^t dt_1 \phi_0(t_1) J_{m_0}[a(t_1)] \exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega+i\hbar\delta)(t-t_1)\right]$$

$$\simeq J_{m_0}[a(t)] \int_0^t dt_1 \phi_0(t_1) \exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega+i\hbar\delta)(t-t_1)\right]$$

$$=\frac{1}{2i\pi} J_{m_0}[a(t)] \int_{\gamma-i\infty}^{\gamma+i\infty} d\lambda\, \phi_0(\lambda) \frac{\exp(\lambda t)}{-\lambda+\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega+i\hbar\delta)}$$

$$=\frac{1}{2i\pi} J_{m_0}[a(t)] \int_{\gamma-i\infty}^{\gamma+i\infty} d\lambda\, \phi_0(\lambda) \frac{\exp(\lambda t)}{\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega+i\hbar\delta)}$$

$$\simeq \pi J_{m_0}[a(t)] \phi_0(t) \delta(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega). \tag{18}$$

Here we have made the assumption that

$$|\lambda| \ll \Gamma, \tag{19}$$

i.e., the observation time is much larger than the decay time of the electron density in the system. The definition of the decay rate Γ is given by Eq. (22) below. Thus, the real part of Eq. (18) can be neglected due to the smallness of $|T_{1p}|^2$ and $|T_{2k}|^2$ with respect to the first term terms in Eq. (17). $J_{m_0}[a(t_1)]$ is considered to be a slow varying function with the period Ω, such that

$$\hbar\Omega \ll \Gamma, \tag{20}$$

where $\Gamma/\hbar$ is a characteristic decay rate for $\phi_0(t)$ in Eq. (18) [see Eq. (22) below].

The Schrödinger equation is now simplified. Therefore, we can reduce it to the following form:

$$\hbar\frac{d\phi_0}{dt} = iJ_{m_0}[a(t)] T_{1p} \exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega + i\hbar\delta)\right] - \Gamma(t)\phi_0(t), \tag{21}$$

with

$$\Gamma \equiv \pi J_{m_0}^2[a(t)] \sum_p |T_{1p}|^2 \delta(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega) + \pi \sum_{n=N}^{\infty} J_n^2[a(t)] \sum_k |T_{2k}|^2 \delta(\epsilon_0-\epsilon_k-V_0/2+n\hbar\omega). \tag{22}$$

In the second term representing decay through the right barrier, there are many resonances due to the wide metal zone. These resonances should be summed up in order to correctly determine the decay rate.

The solution of Eq. (22) is sought in the form

$$\phi_0(t) = b(t) \exp\left[-\frac{1}{\hbar}\int_0^t \Gamma(\tau) d\tau\right]. \tag{23}$$

Substituting Eq. (23) into Eq. (21), one obtains

$$\phi_0(t) = iT_{1p} \int_0^t dt_1 J_m[a(t_1)] \exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2 + m_0\hbar\omega)t_1 - \frac{1}{\hbar}\int_{t_1}^t d\tau\,\Gamma(\tau)\right]. \tag{24}$$

The wave function $\phi_k(t)$ can be found from the substitution of $\phi_0(t)$ determined by Eq. (24), into Eq. (9):

$$\phi_k(t) = iT_{1p}T_{2k} \int_0^t dt_1 \int_0^{t_1} dt_2 J_{n_0}[a(t_1)] J_{m_0}[a(t_2)] \times \exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_k-V_0/2+n_0\hbar\omega)t_1 + \frac{i}{\hbar}(\epsilon_0-\epsilon_p + V_0/2+m_0\hbar\omega)t_2 - \frac{1}{\hbar}\int_{t_2}^{t_1} d\tau\,\Gamma(\tau)\right]. \tag{25}$$

Since $E(t)$ is a periodic function, $\Gamma(t)$ is also periodic with the period

$$T = \frac{2\pi}{\Omega},$$

and, therefore, can be expanded into the Fourier series

$$\Gamma(t) = \Gamma_{\mathrm{eff}} + \sum_{k=1}^{\infty} \left[a_k \cos\left(\frac{2\pi kt}{T}\right) + b_k \sin\left(\frac{2\pi kt}{T}\right)\right], \tag{26}$$

where

$$\Gamma_{\mathrm{eff}} \equiv \frac{1}{T}\int_0^T d\tau\,\Gamma(\tau) \tag{27}$$

is the effective decay rate, or the rate averaged over the period. Thus, by making use of Eq. (26), one finds that

$$\exp\left[\frac{1}{\hbar}\int_0^t d\tau\Gamma(\tau)\right]=\exp\left[\frac{1}{\hbar}\int_0^t d\tau[\Gamma(\tau)-\Gamma_{\rm eff}]+\frac{1}{\hbar}\Gamma_{\rm eff}t\right]. \tag{28}$$

According to Eq. (26), the term $\{\exp(1/\hbar)\int_0^t d\tau[\Gamma(\tau)-\Gamma_{\rm eff})]\}$ is now a slow varying periodic function [see condition (20)]. Consequently, it can be placed in front of the integral in Eq. (25) and canceled out with the same function in the integral over t_1. Finally, one obtains the following simple equation for $\phi_k(t)$:

$$\phi_k(t)\simeq iT_{1p}T_{2k}J_{n_0}[a(t)]J_{m_0}[a(t)]\int_0^t dt_1\int_0^{t_1}dt_2 \times\exp\left[\frac{i}{\hbar}(\epsilon_0-\epsilon_k-V_0/2+n_0\hbar\omega)t_1+\frac{i}{\hbar}(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega)t_2-\frac{1}{\hbar}\Gamma_{\rm eff}(t_1-t_2)\right]. \tag{29}$$

III. THE RESONANT TUNNELING CURRENT

Employing Eq. (29), one can derive the time-dependent probability current of the electron density when $t\to\infty$. We define the transition probability as follows:

$$w_{p\to k}\equiv\left.\frac{d|\phi_k(t)|^2}{dt}\right|_{t\to\infty}=\frac{2\pi}{\hbar}|T_{1p}|^2|T_{2k}|^2J_{n_0}^2[a(t)]J_{m_0}^2[a(t)] \times\frac{\delta(\epsilon_k-\epsilon_p-V_0+(n_0-m_0)\hbar\omega)}{(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega)^2+\Gamma_{\rm eff}^2}. \tag{30}$$

The rate of elastic tunneling may be determined from the following resonance series:

$$T(\epsilon_1^{(z)}\to\epsilon_2^{(z)})=\sum_{p,k} w_{p\to k}\delta(\epsilon_p-\epsilon_1^{(z)})\delta(\epsilon_k-\epsilon_2^{(z)}) =\frac{1}{\pi}J_{m_0}^2[a(t)]\Gamma_1\Gamma_2\left[\sum_{n=N}^{\infty}J_n^2[a(t)]\right] \times\frac{\delta[\epsilon_k-\epsilon_p-V_0+(n_0-m_0)\hbar\omega]}{(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega)^2+\Gamma_{\rm eff}^2}, \tag{31}$$

where the decay rates, Γ_1 and Γ_2, stand for

$$\Gamma_1\equiv\pi\sum_p |T_{1p}|^2\delta(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega),$$

$$\Gamma_2\equiv\pi\sum_k |T_{2k}|^2\delta(\epsilon_0-\epsilon_k-V_0/2+n_0\hbar\omega). \tag{32}$$

Here n_0 and m_0 are defined by resonance condition (16), and N is the number of emitted photons from the localized level inside the well to the right lead electronic states. The current can be calculated from the following expression:

$$J=J_{\Rightarrow}-J_{\Leftarrow}=e\int d\epsilon_1 d\epsilon_2\{T(\epsilon_1^{(z)}\to\epsilon_2^{(z)})f(\epsilon_1)[1-f(\epsilon_2)] -T(\epsilon_2^{(z)}\to\epsilon_1^{(z)})f(\epsilon_2)[1-f(\epsilon_1)]\}. \tag{33}$$

The resultant current (33) includes both forward and backward contributions. The transmission coefficient depends only on the energies $\epsilon_{1,2}^{(z)}$ in the direction of tunneling, while $f(\epsilon)$ is determined by the total electronic energy

$$\epsilon=\epsilon^{(z)}+\frac{k_\perp^2}{2m_e^*}. \tag{34}$$

Integrating out in $k_\perp$ space, one obtains the final expression for the tunneling current through the DBS:

$$J=\frac{em_e^*kT}{2\pi^2\hbar^3}\Gamma_1\Gamma_2\int d\epsilon\left[J_{m_0}^2[a(t)]\sum_{n=N}^{\infty}J_n^2[a(t)] \times\frac{\ln\{1+\exp[(\epsilon_F-\epsilon)/kT]\}}{(\epsilon_0-\epsilon_p+V_0/2+m_0\hbar\omega)^2+(\Gamma_{\rm eff}^{m,N})^2} -J_{n_0}^2[a(t)]\sum_{m=M}^{\infty}J_m^2[a(t)] \times\frac{\ln\{1+\exp[(\epsilon_F-\epsilon)/kT]\}}{(\epsilon_0-\epsilon_k-V_0/2+n_0\hbar\omega)^2+(\Gamma_{\rm eff}^{n,M})^2}\right]. \tag{35}$$

The Fermi energy is assumed to be small with respect to the photon energy $\hbar\omega$. As follows from Eq. (35), the time dependence of the current is due to modulation by laser pulses.

For simplicity, we consider a few specific forms of the time-dependent amplitude $a(t)$.

(1) *cw pulses*. We assume that the amplitude of the electric field can be expressed as follows:

$$E(t)=E_0\cos(\Omega t), \tag{36}$$

where

$$\Omega\ll\omega. \tag{37}$$

In Fig. 2 the time-dependent evolution of the current is depicted for two particular values of the laser intensity parameter a_0, where

$$a_0\equiv\frac{\mu_0E_0}{2\hbar\omega}, \tag{38}$$

with $a_0=1.7$ and $a_0=3.8$. For $a_0=3.8$, negative resistance is found. The current is periodically modulated by the laser field, and the structure of the periodic behavior is rather different for different values of a_0. The temperature is chosen to be very low, and the voltage has the resonance value $V_0/\hbar\omega=2.0$. The electron density coherently oscillates between two wells with the period Ω.

In Fig. 3 we have studied the temperature dependence of current evolution. At high temperatures, the value of the current increases by one order of magnitude. The shape of the time-dependent curve, however, remains largely unchanged.

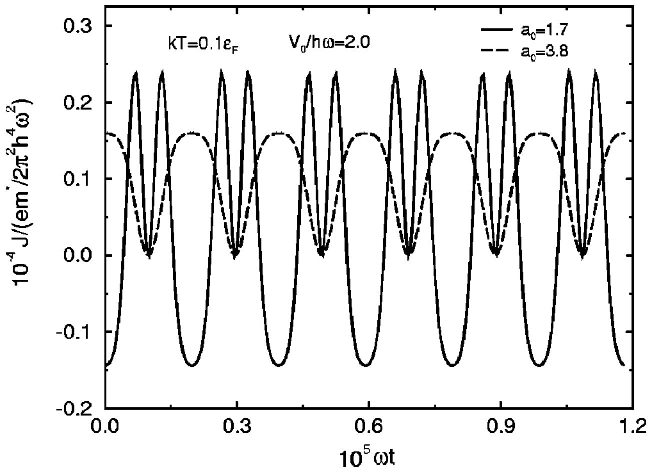

FIG. 2. Time-dependent current modulated by a pulsed electric field with the amplitude $E_0(t)=E_0\cos(\Omega t)$. At $a_0=\mu_0E_0/2\hbar\omega=1.7$ the current is aligned in the positive direction, while at $a_0=3.8$ the current has both positive and negative values.

(2) *Gaussian pulses*. For the Gaussian shape of the periodic pulse, the amplitude, $E(t)$, can be represented in the following form:

$$E(t)=E_0\sum_{-\infty}^{+\infty}\exp[-A(t)], \tag{39}$$

where

$$A(t)=(t-nT_0)^2/\tau_0^2. \tag{40}$$

Here, $E(t)$ is a periodic function with the period T_0. The pulse decay time τ_0 is assumed to be small, i.e.,

$$\tau_0\ll T_0. \tag{41}$$

As demonstrated in Fig. 4, a single input pulse is transformed into two equivalent pulses in the output current. Such transformations performed by the electronic device make a decoding process impossible unless the parameters of the electric field are known.

(3) *Square pulses*. The transformation of square input pulses is the most dramatic. As shown in Fig. 5, the sequence

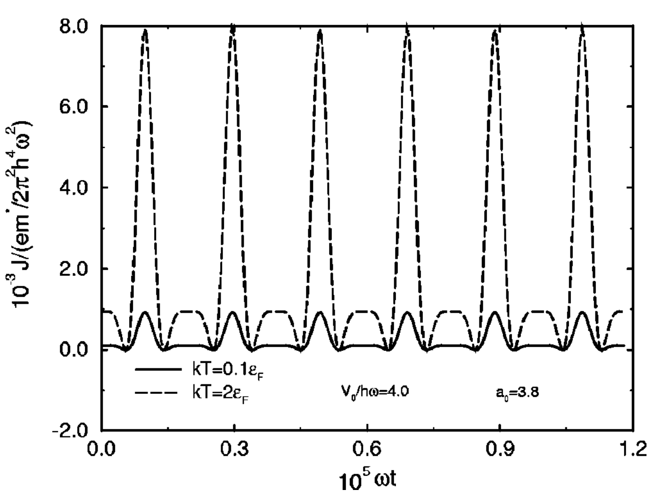

FIG. 3. Temperature dependence of the current. Higher temperatures increase the amplitude of the current oscillations.

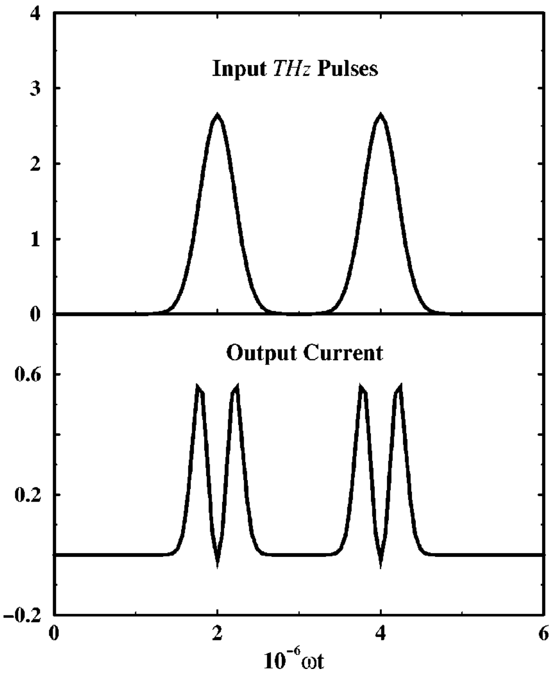

FIG. 4. Transformation of input THz Gaussian pulses into periodic electric current. A single peak in the optical channel corresponds to two equivalent peaks in the electric output. $\omega\tau_0=3\times10^5$, $\omega t_0=2\times10^6$.

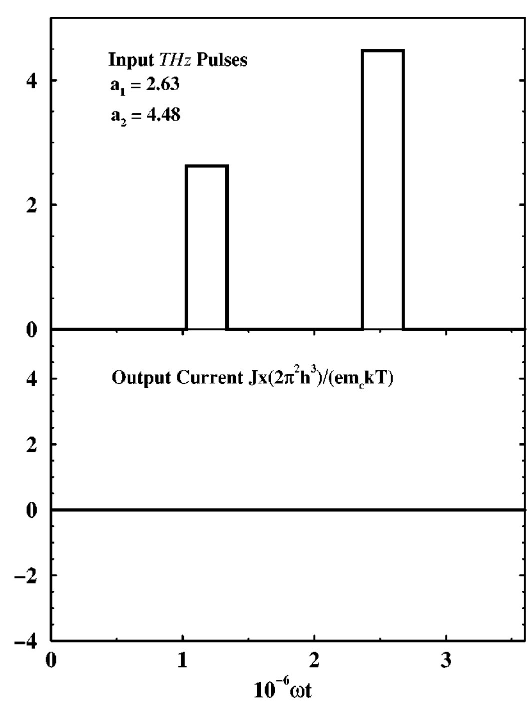

FIG. 5. Sequence of two square optical pulses generates vanishing electric current when $a_1=2.63$ and $a_2=4.48$. $\omega t_0=1\times10^6$.

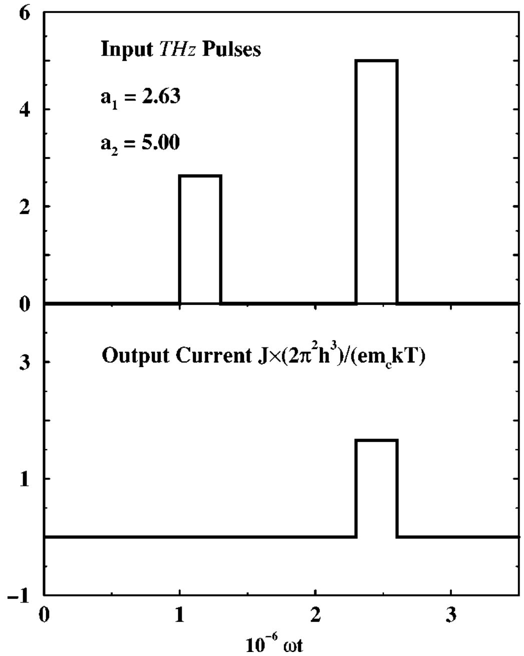

FIG. 6. Sequence of two square input optical pulses is mapped to a single-pulsed output electric current. The values of the intensity parameters for each input pulse are chosen to be $a_1=2.63$ and $a_2=5.00$. $\omega t_0=1\times10^6$.

of two input pulses entirely vanishes in the output electric current when the intensities are correctly tuned ($a_1=2.63, a_2=4.48$). Thus, no output information comes through into electric circuit, and decoding of the input is impossible to achieve.

The quantum well can also be irradiated by a field with two uneven pulses shown in Fig. 6 ($a_1=2.63, a_2=5.00$), resulting in the disappearance of one pulse in the output electric circuit. The output information is, therefore, incomplete, and decoding is impossible as in the previous case. Thus, by choosing the proper intensity of input optical information, one can control the output pulses in such a way as to encode information.

IV. CONCLUSIONS

We have calculated a tunneling current in double-barrier structures irradiated by a strong, slow-varying pulsed THz electric field. We have found that the electronic device essentially changes input signals in the output electric circuit. In particular, depending on the details of the pulse sequence, the current exhibits the following properties.

(a) A current flowing in the direction opposite to applied dc bias, i.e., absolute negative resistance.

(b) The frequency of the current follows the frequency of the optical pulses (see Fig. 2).

(c) The modulated current strongly depends on temperature: the higher the temperature, the larger the current (Fig. 3). The temperature cannot be too high, however, since then the electron-phonon interaction should be taken into account.

(d) At some particular values of the field intensity the input single Gaussian pulse transforms into two even electrical pulses (see Fig. 4). Thus, decoding becomes difficult.

(e) Two nonequivalent THz square pulses with appropriately chosen amplitudes generate null response in the output current (see Fig. 5). Consequently, decoding becomes impossible.

(f) A sequence of two uneven optical pulses with appropriately chosen amplitudes transforms into a single pulse in the electrical circuit (Fig. 6). Hence, one half of the information is lost, and, again, decoding is impossible.

Thus, a double-well structure can be used as a highly nonlinear optoelectronic device. The properties (d), (e), and (f) can be employed to encode input optical information. Decoding becomes impossible if a receiver is not provided with the information about the input parameters. At some values of the electric field, the information is entirely or partially lost during the transmittance.

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