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Resonant tunneling and bistability in a double barrier structure under an intense terahertz laser

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By using exact wave functions of an electron in a terahertz laser field, we calculated the electron resonant tunneling through a double barrier structure. It is found that the laser field has two effects on the current voltage characteristics. First, it introduces additional tunneling states through the structure due to multiple photon processes (emission, absorption, scattering). Second, it reduces the width of the bistable region. At high field strength and low frequencies, the bistability can be completely removed. This can provide a method by which to tune the bistable region in a double barrier structure. © *2001 American Institute of Physics.* [DOI: 10.1063/1.1381033]

In recent years, there has been a rapid development of high-power, long wavelength, tunable laser sources such as free electron lasers (FELs). These radiation sources can provide linearly polarized laser radiation in the terahertz regime[1–8] which have been applied to experimental investigation of nonlinear transport and optical properties in electron gases such as low dimensional systems. Many interesting terahertz phenomena have been investigated, including resonant absorption,[1] the photon enhanced hot-electron effect,[4] THz photon-induced impact ionization,[5] the longitudinal optical (LO)-phonon bottleneck effect,[6] THz photon assisted tunneling,[7] THz cyclotron resonance[8] and dielectric response.[9,10]

In this letter, we study the effect of THz radiation on the electron resonant tunneling through a double barrier resonant tunneling structure (DBRTS). We consider a situation where the terahertz radiation is present in the resonant well and the collector while there is no radiation on the emitter side. The assumption of no radiation on the emitter side is realistic experimentally. We consider a situation where the radiation is incident from the collector side and there is a sufficient amount of free carriers in the collector and the resonant well. In this case this is strong photon absorption in the well and the collector. The number of photons that reach the emitter is thus very small. The resonant level in the well will split into a series of photon side bands. For electron tunneling into the well from the emitter, its energy must be in resonance with one of the photon side bands in the well. There is no additional condition governing the tunneling out of the well to the collector. Transport through the collector is determined by Ohm's law.

Let us consider a two-dimensional electron in the resonant well. Its energy is $\epsilon=(\hbar k)^2/(2m^*)+\epsilon_c$, where ϵ_c is the quasibound energy level in the well. Its state is given by $\psi(\mathbf{r},z)=e^{i\mathbf{k}\cdot\mathbf{r}}\xi(z)$, where $\mathbf{k}$ is the two-dimensional wave vector on the x-y plane and $\xi(z)$ is the localized wave function in the z direction. Here m^* is the effective mass of the electron.

In the presence of an intense laser, the system is strongly coupled to the photon field. Let us choose the vector potential for the laser field to be in the form $\mathbf{A}=(F/\omega)\sin(\omega t)\mathbf{e_y}$. Here F and ω are, respectively, the strength and frequency of the laser field. The time-dependent wave function can then be obtained in term of a unitary transformation.

$$\Psi(\mathbf{r},t)=U\exp(-i\epsilon t/\hbar)\psi(\mathbf{r},z). \tag{1}$$

The unitary transformations U can be written as

$$U=e^{ir_0k_y(1-\cos(\omega t))}\mathbf{e}^{i\gamma(2\omega t+\sin(2\omega t))}, \tag{2}$$

where $\gamma=(eF)^2/(8m^*\hbar\omega^3)$ and $r_0=(eF)/(m^*\hbar\omega^2)$.

Our model Hamiltonian for tunneling through a DBRTS can now be written as

$$H=\sum_{p,\nu,k}(\epsilon_{p,\nu}+\epsilon_k)a^\dagger_{p,\nu,k}a_{p,\nu,k}+\sum_k(\epsilon_c+\epsilon_k)a^\dagger_{c,k}a_{c,k}$$
$$+\sum_{p,\nu,k}[V_{p,\nu}a^\dagger_{c,k}a_{p,\nu,k}+\text{c.c.}]. \tag{3}$$

Here $\nu=l(r)$ refers to the emitting (collecting) lead with $\epsilon_z(\nu)=\epsilon_{p,\nu}$, p is the momentum component perpendicular to the interface, and c refers to the center well with quasibound energy level ϵ_c. $a^\dagger_\alpha(a_\alpha)$ is the creation (annihilation) operator for an electron in the state $|\alpha\rangle$. The tunneling matrix element $V_{p,\nu}$ between the leads and the quantum well is calculated according the Bardeen prescription[11] with one-dimensional potential which includes both the band discontinuities and the applied bias.

By applying the S-matrix scattering theory, the total probability $T_t(\epsilon_z,k)$ for an electron with energy (ϵ_k,ϵ_z) in the emitter to be transmitted to the collector can be written as a product of the elastic coupling to the two leads and the Fourier transform of a Green's function from the resonant level in the well,[12–14]

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$$T_t(\epsilon_z,k)=\Gamma_l\Gamma_r\int_0^\infty ds\int_0^\infty dt e^{i\epsilon_z(t-s)}$$
$$\times\sum_{k'} e^{i\epsilon_{k'}(t-s)}\theta(t)\theta(s)\langle a_{c,k}(t-s)$$
$$\times a^\dagger_{c,k'}(t)a_{c,k'}(t)a^\dagger_{c,k}(0)\rangle, \qquad (4)$$

Here $\Gamma_{l(r)}$ is the decay rate of the resonant level due to the elastic coupling to the left (right) lead, calculated within the wide band approximation.[12] In. Eq. (4) the angular brackets indicate thermal average. For a center well under THz radiation,

$$a_{c,k}(t)=e^{ir_0k_y(1-\cos(\omega t))}e^{i\gamma(2\omega t+\sin(2\omega t))}c_{c,k}(t),$$

where $c_{c,k}(t)$ is the electron operator in the absence of THz radiation. Now the total probability can be written as

$$T_t(\epsilon_z,k)=\Gamma_l\Gamma_r\int_0^\infty ds\int_0^\infty dt e^{i\epsilon_z(t-s)}$$
$$\times\sum_{k'} e^{ir_0k_y(1-\cos(\omega(t-s)))}e^{i\gamma(2\omega(t-s)+\sin(2\omega(t-s)))}$$
$$\times e^{i\epsilon_{k'}(t-s)}\theta(t)\theta(s)\langle c_{c,k}(t-s)c^\dagger_{c,k'}(t)c_{c,k'}(t)c^\dagger_{c,k}(0)\rangle. \qquad (5)$$

Making use of the generating function of the Bessel function the above result can be written as

$$T_t(\epsilon_z,k)=\Gamma_l\Gamma_r\int_0^\infty ds\int_0^\infty dt\sum_{n,m} i^n e^{i(\epsilon_z-2\gamma\hbar\omega)(t-s)}$$
$$\times e^{ir_0k_y}e^{-i(n-2m)\omega(t-s)}J_n(r_0k_y)J_m(\gamma)$$
$$\times\sum_{k'} e^{i\epsilon_{k'}(t-s)}\theta(t)\theta(s)\langle c_{c,k}(t-s)$$
$$\times c^\dagger_{c,k'}(t)c_{c,k'}(t)c^\dagger_{c,k}(0)\rangle. \qquad (6)$$

The above integration can be carried out to yield following result:

$$T_t(\epsilon_z,k)=\frac{\Gamma_l\Gamma_r}{\Gamma}\sum_{n,m}\frac{i^n e^{ir_0k_y}J_n(r_0k_y)J_m(\gamma)}{i(\epsilon_z-\epsilon_c(n,m))+\Gamma/2}, \qquad (7)$$

where $\Gamma=\Gamma_l+\Gamma_r$ and $\epsilon_c(n,m)=\epsilon_c+2\gamma\hbar\omega+(n-2m)\hbar\omega$.

The current through the structure is now given as

$$I(V)=\frac{e}{\pi\hbar}\int d\epsilon[f_l(\epsilon_z)-f_r(\epsilon_z+eV)]\sum_k |T_t(\epsilon_z,k)|^2$$
$$=\frac{e}{\pi\hbar}\int d\epsilon[f_l(\epsilon_z)-f_r(\epsilon_z+eV)]$$
$$\times\sum_k\sum_{n,n',m,m'}\frac{(-1)^{n'}i^{n+n'}J_n(r_0k_y)J_m(\gamma)}{i(\epsilon_z-\epsilon_c(n,m))+\Gamma/2}$$
$$\times\frac{nJ_{n'}(r_0k_y)J_{m'}(\gamma)}{\Gamma/2-i(\epsilon_z-\epsilon_c(n',m'))}, \qquad (8)$$

where $f(\epsilon)$ is the Fermi–Dirac function for the leads. Multiple photon processes (emission, absorption, scattering) lead to a rather complicated tunneling current. The diagonal terms ($n=n',m=m'$) represent a processes in which incoming electrons tunnel through the structure resonantly through one of the photon sidebands. The off-diagonal terms represent photon induced quantum interference in resonant tunneling. An electron can come in and absorb (or emit) $n-2m$ photons then immediately emit (or absorb) $n'-2m'$ photons before tunneling out.

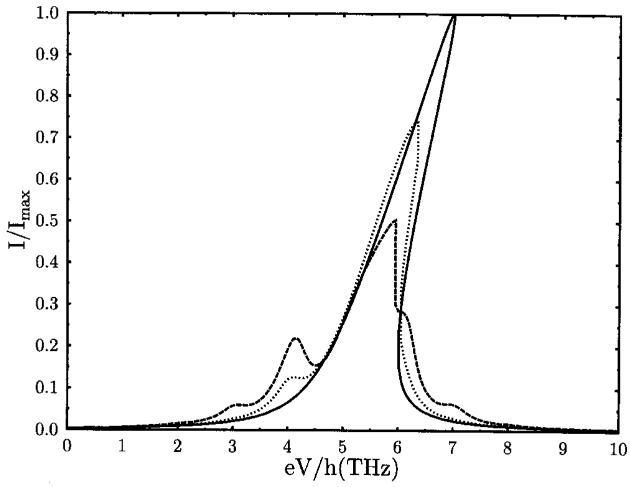

FIG. 1. $I-V$ characteristics in the presence of a THz laser for three difference laser field strength. The laser frequency is 1 THz. The energy difference between the emitter and the center well at zero bias is chosen to 20 meV. The x axis is the applied bias between the emitter and the center well. The Fermi energy of the emitter is 10 meV. The current is in units of $I_{\max}$, the maximum current in the absence of a laser field. The solid line is for $F=0.0$, the dotted line is for $F=10^5$ V/m, and the broken line is for $F=2\times10^5$ V/m.

In the case where the collector barrier is thicker than the emitter barrier, there is considerable charge accumulation in the resonant well. Electron–electron interaction in the well pushes the quasibound energy level upward by an amount δE_{ee}. As a result the current exhibits bistability (or tristability)[15] as a function of the applied bias. In the Hartree–Fock approximation δE_{ee} is directly proportional to the charge density n_e (or current density I, $\delta E_{ee}=aI$). In the present situation where the original resonant level in well splits into photon sidebands due to THz radiation, the electron density in each sideband is proportional to the population of each sideband, $n_e(n,m)=n_e|J_n(r_0k_y)J_m(\gamma)|^2$. The upward shift of a given energy sideband is then given as

$$\delta E_{ee}(n,m)=aI|J_n(r_0k_y)J_m(\gamma)|^2. \qquad (9)$$

This shift is now included in the current given in Eq. (8) and $\epsilon_c(n,m)$ becomes $\epsilon_c(n,m)+\delta E_{ee}(n,m)$. A self-consistent solution of Eq. (8) for I will give rise to tristability in the absence of THz radiation, and such tristability will be affected by the THz radiation. We point out that our analysis does not include a self-consistent solution of Poisson's equation which has the tendency to widen the negative differential resistance region.[16–20] Therefore the effect of space charge due to a self-consistent solution of Poisson's equation is opposite to that due to the laser field. As a consequence, the required strength of the laser field to remove the bistability should be stronger than that calculated here.

In Figs. 1 and 2, we present the calculated tunneling current as a function of applied voltage for various different laser intensities and frequencies. We use the parameters of

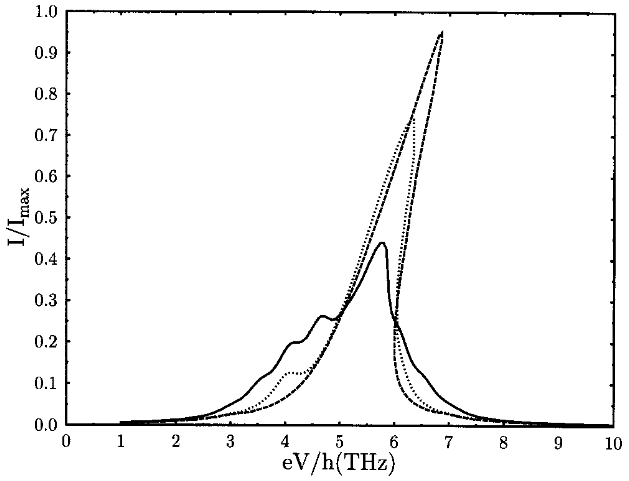

FIG. 2. Same as in Fig. 1, but for different laser frequencies ($f=2\pi\omega$) and at a fixed laser field strength of 10^5 V/m. The solid line is for $f=0.5$ THz, the dotted line is for $f=1.0$ THz, and the broken line is for $f=2.0$ THz.

GaAs based semiconductors, $m^*=0.067m_0$ and $E_F=10$ meV. Figure 1 is the $I-V$ curve at fixed laser frequency of 1 THz. At zero electric field one obtains the expected Z-shaped current, and the width of the tristable region is about 4.1 mV ($eV/h=1$ THz). As the field strength increases to 10^5 V/m, the width of the tristable region decreases to about 1.64 mV ($eV/h=0.4$ THz). At the same time two shoulders appear on both sides of the main resonance due to resonant tunneling through the first order photon sidebands ($m=0,n=\pm 1$). As the field strength increases further ($F=2\times 10^5$ V/m), the width of the tristable region becomes negligible. There are now more side shoulders representing high order photon processes. The shoulders on the low energy side now become satellite resonant peaks, and reflect the strong coupling between electrons and photons. The physical mechanism for the tristability width to decrease can be understood as follows. When the original energy level is split into photon sidebands, electron–electron interaction within each sideband is much smaller compared to that within the original level. The intersideband interaction is also very small in the presence of strong electron–photon coupling.

Figure 2 is the $I-V$ curve at fixed field strength of 10^5 V/m. At high frequencies, ($f=2.0$ THz), the coupling between the electrons and photons is weak. The change of the $I-V$ curve from that of zero field is very small. The reduction of the tristability width is only a few percent. When the laser frequency decreases ($f=1.0$ THz), the coupling increases and the tristability width decreases and additional shoulders develop at lower and higher bias. At low frequency ($f=0.5$ THz), there is very strong coupling between electrons and photons. In this case the tristability width disappears and more satellite peaks and shoulders appear in the tunneling current. The nature of these shoulders and satellite peaks is the same as that described in Fig. 1.

In conclusion, we have shown that, for a DBRTS, THz radiation may be used to produce additional resonance peaks and to tune the width of tristable region in the current–voltage characteristics. The mechanism discussed here can also be useful in detecting THz radiation if the field is of sufficient strength.

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