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LETTER TO THE EDITOR

Temperature-dependent width of current tristability for resonant tunnelling through a double-barrier structure

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Abstract. When electrons tunnel through a double-barrier structure, there exists a region of bistablity (or tristability) in the current–voltage characteristics due to the dynamical charge feedback effect in the resonant well. We propose a mechanism of acoustic-phonon-assisted tunnelling to explain the experimentally observed non-monotonic behaviour of the width of the tristable region. It is found that this width has a single maximum as experimentally observed. Furthermore, contrary to the previous proposal, our result shows that electron thermal activation is not the dominant mechanism in controlling this temperature-dependent width.

The high speed and novel transport properties of the double-barrier resonant tunnelling structure (DBRTS) make it a promising candidate for a new generation of electronic devices. Considerable progress has been made in understanding the physics of DBRTS since the seminal work of Tsu and Esaki [1]. Of particular relevance to the work reported here is the prediction and observation of negative differential resistance (NDR) [2] and intrinsic apparent bistability [3] in the electrical characteristics ($I(V)$ curves) of DBRTSs. Apparent bistability [4, 5] can arise when charge build-up in the central quantum well raises the resonant level in the well (due to the electron–electron interaction) and the resulting potential drop across the collector barrier tips the bell-shaped resonance curve over to the right, giving a Z-shaped (tristable) characteristic, and is particularly pronounced in devices with asymmetric barriers. The width of this tristable region increases with temperature at low temperature and this is usually understood as due to the increased phase space during the tunnelling process [6, 7]. At high temperature, this width decreases monotonically with temperature and this behaviour still remains unexplained. In figure 1, typical experimental results are shown [6]. In what follows, we shall see that such temperature dependence at high temperature can be fully explained with the inclusion of electron–acoustic phonon coupling, i.e., emission and absorption of acoustic phonons by the tunnelling electrons in the resonant well. Our result also strongly indicates that even at low temperature, electron thermal activation (phase space argument) is not the dominant mechanism responsible for the fast increase of the tristability width.

It is well known that electron–LO phonon interactions can give rise to resonant satellites [8]. However, the effect of LO phonons on the temperature dependence of the tristability is not very significant due to the large energy separation between the main resonance and the LO phonon satellite. In this paper we present a calculation of resonant tunnelling through a DBRTS in the presence of electron–acoustic phonon coupling. The theoretical model employed here to treat the electron–acoustic phonon coupling is very similar to those used in the problems of electron–LO phonon coupling [9–11] and deep-electron x-ray emission [12]. Here the acoustic phonon field is expressed in terms of

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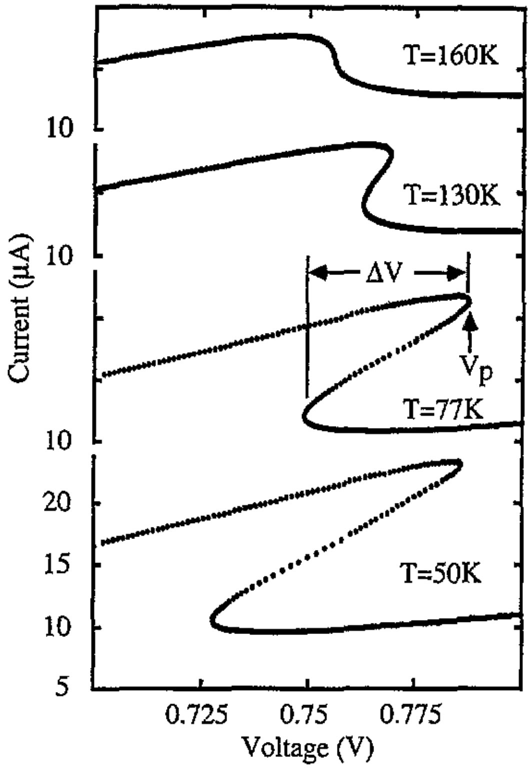

Figure 1. Experimental $I(V)$ curves for a DBRTS, taken from [7], showing the onset of tristability as the temperature is lowered.

an equivalent boson field and electron–acoustic phonon coupling is treated within the deformation potential formalism. We calculate the total energy-dependent transmission probability $T(\epsilon)$, for an electron with energy ϵ incident from the emitter onto the resonant level where it interacts with an acoustic phonon and is then transmitted to the collector. Due to the electron–acoustic phonon coupling and soft dispersion of acoustic phonons, the resulting $T(\epsilon)$ exhibits a shoulder or a weak peak on the high-bias side of the main resonance peak, representing the inelastic-scattering-assisted resonant tunnelling.

Let the total energy $\epsilon = \epsilon_k + \epsilon_z$ where $\epsilon_k = \hbar^2 k^2/2m$ (k is the momentum along the plane) and ϵ_z is the energy perpendicular to the plane. The phonon energy is $\hbar\omega_Q$ where Q is a three-dimensional wavevector ($Q = q, q_z$). Our model Hamiltonian consists of three terms, $H = H_e + H_p + H_I$, where H_e is the Hamiltonian for the electrons including the elastic coupling to the tunnelling barriers,

$$H_e = \sum_{p,\nu,k} (\epsilon_{p,\nu} + \epsilon_k) a^\dagger_{p,\nu,k} a_{p,\nu,k} + \sum_k (\epsilon_c + \epsilon_k) a^\dagger_{c,k} a_{c,k} + \sum_{p,\nu,k} \left[V_{p,\nu} a^\dagger_{c,k} a_{p,\nu,k} + \mathrm{cc} \right]. \quad (1)$$

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 centre well with quasibound energy level ϵ_c. The tunnelling matrix element $V_{p,\nu}$ between the leads and the quantum well is calculated according to Bardeen's prescription [13] with a one-dimensional potential which includes both the band discontinuities and the applied bias. The term H_p is a boson-like field Hamiltonian representing the acoustic phonon,

$$H_p = \sum_Q \hbar\omega_Q b_Q^\dagger b_Q \tag{2}$$

where ω_Q is the phonon frequency and only its Debye approximation will be used below. The electron–phonon coupling H_I can be written as

$$H_I = \sum_{k,Q} M_Q(b_Q^\dagger + b_Q)a_{c,k+q}^\dagger a_{c,k} \tag{3}$$

which only allows momentum transfer along the plane. In the above expression, we use M_Q to denote the electron–phonon coupling strength [11], $M_Q = D(\hbar/2\rho v\omega_Q)^{1/2}Q$, where D is the deformation potential parameter, ρ the lattice density, and v the volume.

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 function from the resonant level in the well [10, 11],

$$T_t(\epsilon_z, k) = \Gamma_l\Gamma_r \int_0^\infty \mathrm{d}s \int_0^\infty \mathrm{d}t\, \mathrm{e}^{\mathrm{i}\epsilon_z(t-s)} \sum_{k'} \mathrm{e}^{\mathrm{i}\epsilon_{k'}(t-s)}\theta(t)\theta(s)$$
$$\times\langle a_{c,k}(t-s)a_{c,k'}^\dagger(t)a_{c,k'}(t)a_{c,k}^\dagger(0)\rangle. \tag{4}$$

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 [10]. In equation (4) the angular brackets indicate the thermal average. The four-operator correlator can be calculated using the cumulant expansion method [11, 13] which includes the high-order vertex corrections. We obtain

$$T_t(\epsilon_z, k) = \frac{2\Gamma_l\Gamma_r}{\Gamma} \int_0^\infty \mathrm{d}t\, \exp\left[\mathrm{i}\Omega_k t - \frac{1}{2}\Gamma t + \sum_Q F_Q\right] \tag{5}$$

where $\Gamma = \Gamma_l + \Gamma_r$, $\Omega_k = (\epsilon_z - \epsilon_c + \lambda_k)$ and

$$\lambda_k = \sum_Q \frac{|M_Q|^2}{\omega_Q^2 - \Delta^2}\left[\omega_Q + (1 + 2N_Q)\Delta\right]. \tag{6}$$

Here the Q summation is restricted to the Debye wavevector. In equation (6), the phonon population $N_Q = (\exp(\beta\omega_Q) - 1)^{-1}$ and $\Delta = \epsilon_k - \epsilon_{k-q}$. The function F_Q in equation (5) is given as

$$F_Q = -\frac{|M_Q|^2}{(\omega_Q^2 - \Delta^2)^2}\left[(\omega_Q^2 + \Delta^2)(2N_Q + 1) + 2\omega_Q\Delta\right]$$
$$+\frac{|M_Q|^2(\omega_Q^2 + \Delta^2)}{(\omega_Q^2 - \Delta^2)^2}\left((2N_Q + 1)\cos(\omega_Q t) - \mathrm{i}\sin(\omega_Q t)\right)\mathrm{e}^{\mathrm{i}\Delta t}$$
$$-\frac{2|M_Q|^2\omega_Q\Delta}{(\omega_Q^2 - \Delta^2)^2}\left(\mathrm{i}(2N_Q + 1)\sin(\omega_Q t) - \cos(\omega_Q t)\right)\mathrm{e}^{\mathrm{i}\Delta t}. \tag{7}$$

It should be mentioned here that the complex structure of F_Q and $T_t(\epsilon_z, k)$ does not alter the sum rule for the total tunnelling rate. Integrating over ϵ_z introducing a factor $\delta(t)$ ensures the sum rule because $F_Q(t = 0) = 0$, regardless of the k value.

In the following we shall neglect recoil or coherence breaking processes, i.e. we neglect Δ in the tunnelling rate to restrict ourselves to a one-dimensional tunnelling model [9]. We make this approximation to simplify the calculations but also because for larger values of k

the electron–acoustic phonon coupling is greatly reduced if in-plane energy loss is included. This arises from the condition that the energy of the phonon emitted must be greater than or at least equal to the energy lost in the plane.

The current through a DBRTS is given as

$$I(V) = \frac{e}{\pi\hbar} \int d\epsilon \, T(\epsilon)[f_l(\epsilon) - f_r(\epsilon + eV)] \tag{8}$$

where $f(\epsilon)$ is the Fermi–Dirac distribution function for the leads. The case where Δ is set to zero is equivalent to $k = 0$ and thus there is no k dependence of $T(\epsilon)$. Thus the integration can be done analytically, with the distribution function part leading to a temperature-dependent logarithmic factor $g(T)$. The result for a 2D emitter model is presented in figure 2. The broadened secondary peak on the high-bias side of the main peak is due to the emission of acoustic phonons by the tunnelling electrons in the resonant site. Compared to the case of electron–LO phonon coupling, here the acoustic phonon satellites are much broadened because of the strong Q-dependent dispersion. They are also sensitive to the temperature variation. The phonon satellites increase as temperature increases, while the main resonance exhibits a maximum somewhere close to 90 K. The main resonance increases by 20% at around 100 K and reduces by 20% at around 250 K, compared to the zero-temperature case. Also as temperature increases, the phonon population increases and therefore absorption processes start to contribute. This absorption-assisted tunnelling can be seen clearly as the current at bias lower than the main resonance increases with the temperature. Unlike emission processes, the tunnelling current due to absorption only has a weak shoulder.

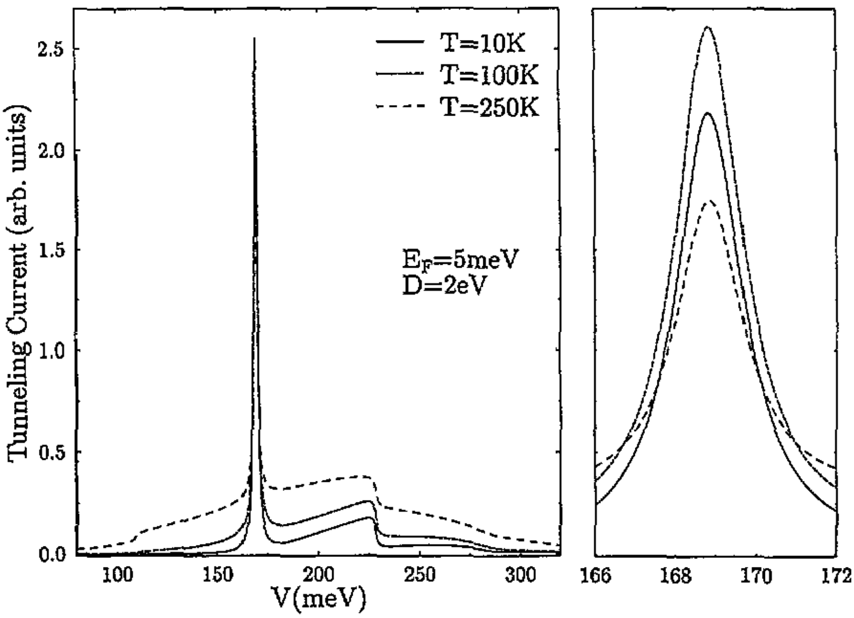

Figure 2. Plot of tunnelling current as a function of applied bias for several different temperatures. The right panel shows how the height of the main resonance varies with temperature.

The tristability width is determined by both the phonon-assisted tunnelling and the electron thermal activation. If electron–phonon interaction is absent, the tunnelling current is given as

$$I = I_{max}\frac{\Gamma^2}{(e(V-V_{th})-\alpha I)^2+\Gamma^2}g(T) \tag{9}$$

where V_{th} is the difference between the 2D energy level in the emitter and that in the well at $V = 0$; αI is the dynamical level shift in the well due to electron–electron interaction. The temperature-dependent statistical factor is given as

$$g(T) = \frac{T}{E_F}\ln\left[\frac{1+\exp[(E_F-E_e)/T]}{1+\exp[(E_F-E_e-eV)/T]}\right]$$

with E_e the energy level of the emitter. This algebraic equation can be solved selfconsistently. The resulting tristability width is a monotonic function of T with exponentially slow increase for $T < T_F/2$ and a roughly linear T dependence for $T > T_F/2$ where T_F is the Fermi temperature of the emitter. Therefore it is clear that a simple activation picture cannot explain the observed temperature-dependent width. On the other hand, if we completely neglect the thermal activation, the phonon-limited width decreases with temperature monotonically (starting from a finite width at zero temperature).

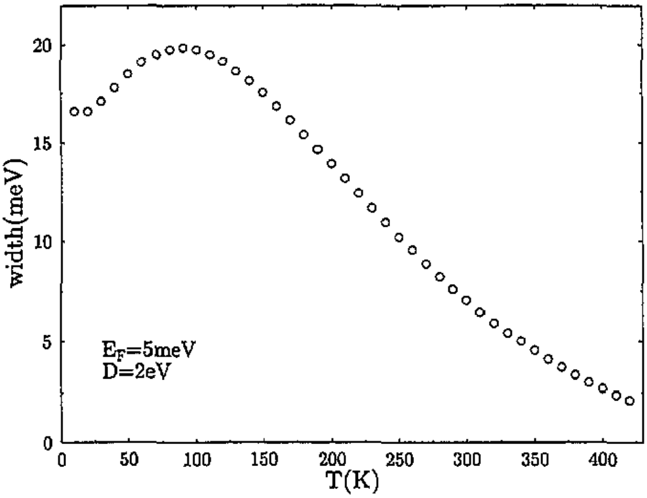

Figure 3. Plot of tristability width, ΔV, against temperature.

We have calculated the tristability width by taking into account both electron activation and acoustic-phonon-assisted tunnelling. The result is plotted in figure 3. It can be seen that the width calculated within this model does indeed have a single maximum at some intermediate temperature (50–100 K) [16] and goes to zero at high temperature. While the present calculation agrees qualitatively well with experiment at intermediate and high temperatures and displays the required single maximum, a discrepancy exists at low temperature. This indicates that the low-temperature width is not controlled by the activation, but rather by another mechanism of stronger temperature dependence. A possible candidate is electron–impurity or electron–disorder elastic scattering, as such elastic-scattering-assisted tunnelling can enhance the maximum tunnelling current. However a full quantitative understanding of the low-temperature width can only be achieved with a calculation of temperature-dependent disorder-assisted tunnelling, which is beyond the scope of the present work. Thus the observed width can be understood as follows. At low temperature, our qualitative understanding is that the width is controlled by both the electron–disorder interaction and thermal activation. A combination of the two mechanisms leads to an increasing ΔV with temperature. As temperature further increases, the phonon-assisted tunnelling starts to dominate both the maximum tunnelling current and the electron tunnelling lifetime in the resonant well. As a result, the tristability width decreases with temperature. The transition from the non-phonon- to the phonon-dominated regime is determined by the temperature at which the width is maximum.

In conclusion, the inelastic resonant tunnelling due to electron–acoustic phonon coupling in a DBRTS is investigated. The temperature-dependent width of current tristability can be calculated. It is found that electron thermal activation plays only a minor role in determining the tristability width.

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