Theory of photo-induced resonant tunneling in heterojunctions

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The theoretical basis for understanding photo-assisted resonant tunneling is generally taken to be the theory by Tien and Gordon [Phys. Rev. **129**, 647 (1963)] in which the effect of a photon field on an energy state is to create sidebands at multiples of the photon energy. An alternative mechanism in which bulk photoabsorption is the photoexcitation mechanism is discussed. The two theories take the same form in the weak-field limit, and an expression is given for the relative size of tunneling currents.

The behavior of resonant tunneling heterojunctions has been the subject of much study since the pioneering work of Esaki and Tsu.[1] This type of heterojunction exhibits useful characteristics for electronic applications, such as oscillators or switching devices, due to its negative differential resistance (NDR) and rapid charge transfer. Here we focus on the behavior of resonant junctions subjected to photon radiation which has been the subject of much interest.[2]

Sollner *et al.* have used resonant tunneling components to observe detection and mixing at frequencies as high as 2.5 THz, oscillations up to 420 GHz at room temperature, and picosecond switching times.[3–7] These junctions have also been used in transistors.[8] These experimental developments have stimulated much theoretical work with a view to understanding the fundamental time scales of a resonant tunneling device.[5,6,9–18]

This paper compares two mechanisms that may explain the experimental results of Sollner *et al.* Theoretical work on resonant tunneling junctions in the presence of a photon field has generally been carried out within the framework first introduced by Tien and Gordon (TG) in 1963.[19] They suggested that in the presence of a photon field an energy level E develops sidebands at $E\pm n\hbar\omega$ where n is an integer. This allows a tunneling current not only from the state at energy E but also at energies corresponding to the sidebands. The model admits a simple analytical solution that will be reviewed below. This picture is in sharp contrast to the way in which photon absorption in solids takes place: as an excitation of an electron from one state to a different one that is separated from it in energy by $\hbar\omega$. We suggest that a second mechanism in which an electron is photoexcited to another state and then tunnels through the barrier may be the appropriate model for some of the experiments on tunneling junctions. In this paper we present the consequences of such a model and compare them with the results of TG.

TG assumed that the effect of the photon field is to produce a spatially constant potential $v\cos(\omega t)$, on one side of the junction. The other side of the junction is taken to be at zero potential. If the wave function in the absence of photons is $f(\mathbf{r})e^{-iEt/\hbar}$ for a state of energy E then in the presence of the potential the wave function above is multiplied by $\Sigma_{-\infty}^{\infty}J_n(\alpha)e^{-in\omega t}$ where $J_n(\alpha)$ is a Bessel function of order n and

$$\alpha=ev/\hbar\omega. \tag{1}$$

It follows that if the tunneling current in the absence of photons ($\alpha=0$) is $j_0(V)$ where V is the dc bias voltage across the junction then in the presence of a photon field the current is

$$j(V,\omega)=\sum_{n=-\infty}^{\infty}J_n^2(\alpha)j_0\left(V+\frac{n\hbar\omega}{e}\right), \tag{2}$$

as shown by TG.[19] The photon field produces sidebands at energies $E\pm n\hbar\omega$ that have a time-averaged probability of occupation $J_n^2(\alpha)$, hence the weighting of $j_0(V+n\hbar\omega/e)$ in Eq. (2). This equation is very appealing; it takes multiphoton processes into account and expresses the current $j(V,\omega)$ in terms of j_0, the current in the absence of the external field.

In the limit $\alpha\ll 1$, Eq. (2) yields

$$j(V,\omega)\simeq j_0(V)+(\alpha^2/4)[j_0(V+\hbar\omega/e)+j_0(V-\hbar\omega/e)-2j_0(V)]. \tag{3}$$

The first term is the current in the absence of photons (elastic current) and the remaining terms are due to single-photon absorption or emission. The term $-2j_0(V)$ in the square bracket is the reduction of the elastic current, $j_0(V)$, by a factor $-\alpha^2/2$ due to the photon-induced inelastic currents, $j_0(V\pm\hbar\omega/e)$, with probability $\alpha^2/4$.

We now very briefly discuss some additional consequences of the TG theory that follow from Eq. (2). In the limit $\alpha\gg 1$, Eq. (2) reduces to the semiclassical result,

$$j(V,\omega)\simeq\overline{j_0(V+v\cos\omega t)}, \tag{4}$$

where the bar indicates a time average. Equation (4) is obtained by using the asymptotic expression for the envelope of $J_n(\alpha)$ in Eq. (2): $J_n^2(\alpha)\propto 1/[\pi\sqrt{\alpha^2-n^2}]$ for $n\ll\alpha$. Expanding the right-hand side of Eq. (4) in a Taylor series gives to lowest order

$$j(V,\omega)-j_0(V)\approx\frac{v^2}{4}\frac{\partial^2 j_0(V)}{\partial V^2}, \tag{5}$$

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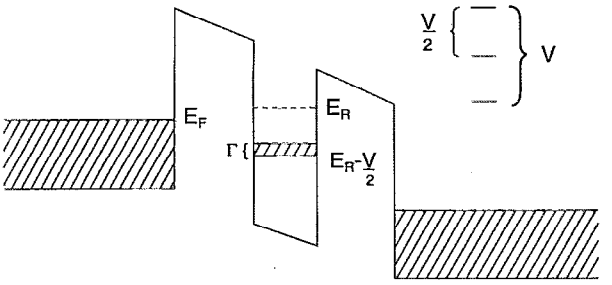

FIG. 1. Energetics for the double-barrier quantum well used in the experiment by Sollner *et al.* (see Ref. 3). The intermediate state is at 78 meV (dashed) with a broadening Γ and a barrier height of 250 meV. Based on the doping of the electrodes the Fermi energy E_F is 53 meV and the plasma frequency ω_p is 39 meV, using an effective mass of $m^* \approx 0.07$ m and a background dielectric function of 13.

which is the small signal result. Equation (5) is valid for $\alpha \gg 1$ and $ev \ll \Gamma$, in other words $\hbar\omega \ll ev \ll \Gamma$, where Γ is the energy scale over which j_0 varies, i.e., $|\partial j_0/\partial V| \simeq e j_0/\Gamma$. Equation (3), which is valid in the limit $\alpha \ll 1$, also leads directly to Eq. (5) if $\hbar\omega \ll \Gamma$. In fact, Eq. (5) is a consequence of Eq. (2) in the limit $\hbar\omega, ev \ll \Gamma$ independent of the relative size of $\hbar\omega$ and ev. This can be shown by observing that Eq. (1) can be written in the operator form

$$j(V,\omega) = J_0\left[2i\alpha \sinh\left(\frac{\hbar\omega}{2e}\frac{\partial}{\partial V}\right)\right] j_0(V), \tag{6}$$

which yields Eq. (5) if $\hbar\omega, ev \ll \Gamma$.

TG derived their results for the special case of a junction in which the tunneling probability is independent of the applied voltage. For resonant tunneling in a double barrier this assumption does not apply because electrons can tunnel only if they are near the resonant energy. To a good approximation, in the presence of a bias voltage V, the resonant energy E_R, in a symmetric junction is shifted by $eV/2$ to $E_R - eV/2$ (see Fig. 1). The presence of a photon field allows electrons with energy $E_R - eV/2 \mp n\hbar\omega$ to tunnel and this current is given by $j_0(V \pm 2n\hbar\omega/e)$. The model of TG can be extended to treat this situation by characterizing the inelastic tunneling by the effective voltage $V \pm 2n\hbar\omega/e$. Thus in Eq. (1), $j_0(V \pm n\hbar\omega/e)$ is replaced by $j_0(V \pm 2n\hbar\omega/e)$ and the energy scale Γ over which j_0 varies corresponds to the width of the resonant state.

Our theory is based on a very different physical process from that described in the TG theory. The TG theory assumes constant potentials in the electrodes, with a jump over the barrier, which is based on the complete screening of the photon field in the electrodes at low frequencies. This spatially constant potential is responsible for the simple form (Bessel functions) of the photon sidebands. It also has the consequence that there is no photon absorption in the absence of tunneling, i.e., it does not require bulk absorption. On the other hand, our theory focuses on the inhomogeneities in the electrodes that lead to bulk photon absorption (Drude absorption). Our formulation invokes a two-step process; bulk-like photoabsorption followed by tunneling through the barrier.

This two-step model yields an expression for the tunneling current that is of the same form as Eq. (3),[18]

$$j_B(V,\omega) = j_0(V) + P[\,j_0(V+2\hbar\omega/e) + j_0(V-2\hbar\omega/e) - 2j_0(V)\,], \tag{7}$$

with the coupling strength

$$P = \left(\frac{\epsilon_0 \epsilon_2 E_e^2}{\hbar}\right)(l)\left(\frac{1}{n v_F}\right). \tag{8}$$

This expression depends mainly on the properties of the electrodes; ϵ_2 is the imaginary part of the dielectric function, l is the electron mean free path, n is the doping density of electrons, v_F is the Fermi velocity, E_e is the electric field in the electrodes, and ϵ_0 is the permittivity of vacuum. The field inside these devices is a complicated function of the externally applied field. The first two factors, delineated by parentheses, combine to give the flux incident on the junction due to the photon absorption; the first factor is the number of photons absorbed per unit volume and time in the electrodes, and the second is the mean free path of the excited electrons. On average, only those electrons excited within a mean free path reach the junction without scattering. The third factor is one over the incident unexcited flux making P the ratio of the excited flux to the unexcited flux. The predominant absorption mechanism in Sollner's experiment is Drude absorption.[20] In this case the imaginary part of the dielectric constant is approximately $\epsilon_2 \approx (1/\omega\tau)(\omega_p/\omega)^2$, where ω_p is the plasma frequency. Since $l = v_F\tau$, P is independent of the scattering rate in the leads. This independence arises because the photon absorption is proportional to the scattering rate, but the fraction of the excited electrons that reaches the junction is inversely proportional to the scattering rate.

Equation (7) was derived with the assumption that the probability of light absorption is the same for all Fermi sea electrons that contribute to the photoabsorption. This should be a reasonable assumption for the case of Drude absorption. The junction current is largest when the bias voltage causes the resonant level to be degenerate with the conduction-band states of the left-hand-side electrode (see Fig. 1). In this case, the first two terms in parentheses of Eq. (7) correspond to an electron tunneling through the resonant level and leaving a conduction-band hole which is then filled by another electron through the process of photon emission or excitation.

The ratio of $\alpha^2/4$, appearing in Eq. (3) of the TG theory, to P in Eq. (8) is seen to be

$$r \equiv \alpha^2/4P = (d/d_0)^2, \tag{9}$$

where d is the barrier thickness, and $d_0 \equiv 2\sqrt{\hbar/m^*\omega}$ for an electromagnetic field with frequency ω and an electron effective mass of m^*. To obtain this result we have used $v = dE_b$, where E_b is the field in the barrier and v is the voltage drop across the junction. We wish to emphasize that the electrode potential assumed by TG to be $v\cos(\omega t)$ is only an approximation since the electric field in the elec-

trodes is not perfectly screened due to the time dependence of the fields. The penetration depth of the field in the electrodes is $c/(\omega \operatorname{Im}\sqrt{\epsilon_e})$, where ϵ_e is the dielectric function of the electrodes. The penetration depth is typically much larger than the electron mean free path and is $\approx 10^5$ Å in the Sollner experiment. Thus, the existence of the TG effect implies that the field penetrates the electrodes and causes bulk absorption, leading to the effect we have treated.

In the case that the field is normal to the barrier, the field in the electrodes E_e is related to the field in the barrier E_b by $\epsilon_e E_e = \epsilon_b E_b$. The dielectric function of the electrodes can be approximated[20] by adding to the static dielectric function of intrinsic GaAs (ϵ_e^0) a free-electron contribution yielding

$$\epsilon_e(\omega) = \epsilon_e^0 - \frac{\omega_p^2}{\omega(\omega + i/\tau)}, \tag{10}$$

where ω_p is the plasma frequency corresponding to the doping and τ is the relaxation time. For the case that Sollner studied, $\epsilon_e(\omega) \approx -3 + 6i$ and $\epsilon_b \approx 13$ are of the same magnitude, as assumed in deriving Eq. (9). A similar conclusion holds for the case of the fields parallel to the barrier, for which $E_e = E_b$. In that situation the TG expression for α is modified from Eq. (1) by a multiplicative factor $2v_F/\omega d$, where v_F is the Fermi velocity.

In the case of the Sollner experiment with $\omega = 10$ meV and $m^* \approx 0.07$ m we obtain $d_0 \approx 100$ Å. The barrier thickness is also approximately 100 Å yielding $r \simeq 1$. Until more detailed tests are made we cannot say which physical picture is the correct one to describe the Sollner experiment. It is clear that the result for r can be quite different depending on the experimental conditions. We would therefore encourage experimental work on the detailed dependence on specific material parameters and characteristic quantities such as the mean free path and the electron density (doping) in the heterojunction. Such work could distinguish the two different physical mechanisms, since they occur simultaneously.

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