Finite-Temperature Fermi-Edge Singularity in Tunneling Studied Using Random Telegraph Signals

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We show that random telegraph signals in metal-oxide-silicon transistors at millikelvin temperatures provide a powerful means of investigating tunneling between a two-dimensional electron gas and a single defect state. The tunneling rate shows a peak when the defect level lines up with the Fermi energy, in excellent agreement with theory of the Fermi-edge singularity at finite temperature. This theory also indicates that defect levels are the origin of the dissipative two-state systems observed previously in similar devices.

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In small electrical devices, noise signals are often seen which reflect the transitions of a single atom or electron between two or more metastable states. These "random telegraph signals" (RTSs) [1], where the conductance jumps randomly in time between certain discrete values, are being increasingly exploited as a means of investigating a diverse range of tunneling phenomena, such as dissipative tunneling of two-state systems [2], electromigration [3], hopping conduction [4], and tunneling between quantum Hall edge channels [5]. In the present work we use them for the first time to study the dynamics of electrons tunneling between a defect state and a two-dimensional electron gas (2DEG) at millikelvin temperatures. We find that a noninteracting electron picture cannot explain the behavior of the system. This is not very surprising, because the interaction phenomena of the Coulomb blockade [6] and the Kondo effect [7] are known to strongly influence the tunnel conductance through a single defect level [8,9] or a quantum dot [10]. However, what is surprising is that the dominant interaction effect in the present system is another, to which attention has only recently been drawn [11]. It is the interaction of the 2D gas electrons with the defect potential which produces a peak in the tunneling rates near the Fermi level at low temperature T. At $T = 0$ this peak becomes a Fermi-edge singularity, whose origin is the same as that of the x-ray absorption edge singularity in metals [12].

This singularity has already been invoked to explain sharp peaks observed in the current-voltage characteristics of resonant tunneling diodes [9], but the interpretation in that situation was hampered by the lack of equilibrium due to the large voltage bias, and by the simultaneous presence of other anomalous structures in the device characteristics. In our RTS experiments we can measure the tunneling rates for an isolated defect directly and in thermal equilibrium. We find excellent agreement with the theory for the finite temperature generalization of the Fermi-edge singularity, which has not been tested experimentally before. Using the same theory we are then able to attribute the weakly coupled dissipative two-state systems observed previously [13] in similar devices to an electron tunneling between a defect level and a bound state of the defect potential. We also find that the effects of a magnetic field B on an RTS are fully consistent with the electron-trapping defect scenario.

The measurements were made on two-terminal Si metal-oxide-semiconductor field-effect transistors (MOSFETs) with highly doped contact regions, oxide thickness $d_{\rm ox} = 240$ Å, and channel dimensions 0.6 or 0.8 μm. Care was taken not to stress the devices electrically, and the threshold gate voltage was about 2.0 V in a dilution refrigerator at 100 mK. The resistance was sampled at up to 5 kHz using a standard constant-current lock-in technique with a Brookdeal 5004 ultralow noise voltage preamplifier. The low-temperature peak mobility was around 0.2 m^2 V^{-1} s^{-1}, corresponding to a transport scattering length $l \sim 300$ Å. This rather high disorder leads to quantum interference effects (see later) which fortuitously make the RTSs big enough to allow the use of signal levels not significantly larger than kT/e (to avoid resistive heating) down to 100 mK. The mean time spent in each resistance state of an RTS under fixed conditions was found by averaging over several hundred transitions, giving a standard error of a few percent.

Previous work at $T \geq 4.2$ K has shown that most RTSs in n-channel MOSFETs result from defect levels situated in the oxide close to the Si/SiO_2 interface [1]. Figure 1(a) shows a schematic band diagram illustrating the situation at low T, together with a section of a typical RTS. A positive voltage V_g is applied to the gate, sufficient to create a degenerate 2DEG, represented by the shaded area at the interface. E_d is a defect level (indicated as being in the oxide, though its location is not important here), E_F is the Fermi level, and E_0 is the bottom of the lowest 2D subband. We denote the conductance of the device with the defect level empty (state 1) or occupied (state 2) by G_1 or G_2, respectively, and the reciprocals of the mean times spent in the two states by rates γ_1 and γ_2, as indicated in the figure. The ratio of these rates satisfies the detailed balance condition,

$$\gamma_1/\gamma_2 = \exp[-(E_d - E_F)/kT]. \qquad (1)$$

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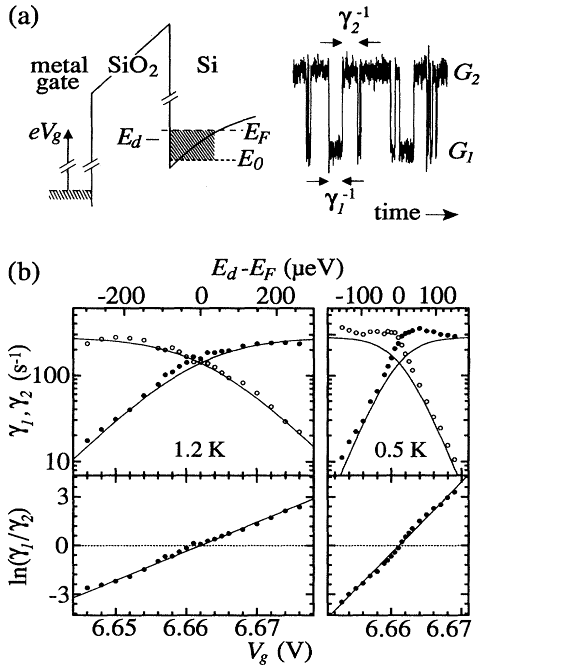

FIG. 1. (a) Schematic band diagram of a MOSFET at low temperature, and a section of a typical RTS seen in such a device. (b) Gate-voltage dependence of the capture rate γ_1 (filled circles) and emission rate γ_2 (empty circles) for RTS 1 at 1.2 and 0.5 K, with $B = 0.28$ T. The device conductance was 2 mS. The solid lines are fits by Eqs. (1)–(3), which neglect interactions.

The lower graphs in Fig. 1(b) show $\ln(\gamma_1/\gamma_2)$ plotted against V_g at two temperatures. The variation is almost linear, implying that

$$E_d - E_F = -\eta e(V_g - V_{g0}), \qquad (2)$$

where η (the "sensitivity") and V_{g0} (the "balance" gate voltage) are constants. The straight solid lines in the figure correspond to $\eta = 0.019$ and $V_{g0} = 6.662$ V. The linearity results from the energy-independent density of states in the inversion layer and the linear sensitivities of E_d and E_0 to the oxide electric field over the relevant range of V_g. Differences in η between individual RTSs can be attributed to different defect locations. The deduced values of $E_d - E_F$ for RTS 1 are plotted along the top axes in Fig. 1(b).

In a noninteracting picture the individual rates, deduced from the golden rule, are given by

$$\gamma_1 = (2\pi/\hbar)D\Delta^2 f(E_d),$$
$$\gamma_2 = (2\pi/\hbar)D\Delta^2[1 - f(E_d)], \qquad (3)$$

where $f(E)$ is the Fermi function, D is the electron density of states, and Δ is the tunneling matrix element. The lines on the upper graphs in Fig. 1(b) are plots of Eqs. (3), where $(2\pi/\hbar)D\Delta^2 = 280\ \mathrm{s}^{-1}$ is the only fitting parameter. At $T = 1.2$ K the measured values of γ_1 and γ_2 follow them rather well; hence the picture of tunneling without interactions seems to suffice. However, when T is reduced to 0.5 K a distinct peak appears in the rates in the region of $E_d = E_F$. To account for this peak, which occurs to some degree for every RTS, we are forced to go beyond the noninteracting picture. Matveev and Larkin [11] have recently pointed out that one should take into account the consequences of the change in the defect potential seen by electrons in the 2DEG when an electron tunnels. Most of the relevant theory was developed in the context of the x-ray absorption edge [12], and at $T = 0$ it predicts a power-law singularity in the transition rate,

$$\gamma_1 \sim \theta(E_F - E_d)(E_F - E_d)^{\alpha - 1}, \qquad (4)$$

where $\theta(E)$ is the unit step function. The singularity arises because the tunneling electron can easily lose a small amount of energy to low-energy electron-hole pairs which are created by the sudden change in the defect potential. The finite temperature generalization of Eq. (4) is [14]

$$\gamma_{1,2} = CT^{\alpha-1}\exp\left(\pm\frac{E_F - E_d}{2kT}\right) \times \frac{|\Gamma[\alpha/2 + i(E_d - E_F)/2\pi kT]|^2}{\Gamma(\alpha)}, \qquad (5)$$

where C is a constant and α is equal to the zero-temperature exponent. Figure 2 shows the transition rates for RTS 2, found in the same device as RTS 1 at lower gate voltage ($V_{g0} = 3.0863$ V and $\eta = 0.018$). The solid lines are the result of fitting by Eq. (5), yielding $\alpha = 0.21 \pm 0.01$ and $C = 10.2 \pm 0.2\ \mathrm{s}^{-1}$ (with T in degrees kelvin). The very good fit at both $T = 145$ and 360 mK is convincing evidence that these measurements directly probe the Fermi-edge singularity at finite temperature. The behavior of other RTSs was completely consistent

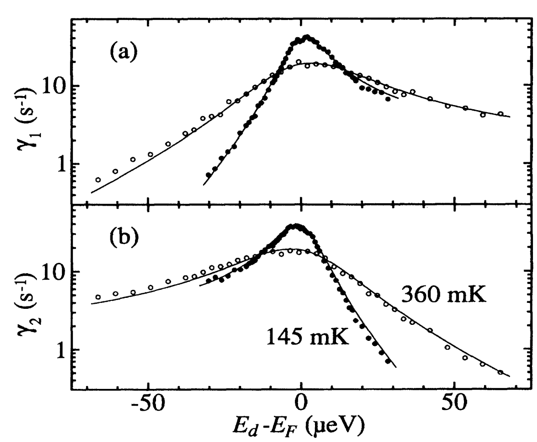

FIG. 2. Energy dependence of (a) capture and (b) emission rates for RTS 2 at 145 mK (filled circles) and 360 mK (empty circles) at $B = 0.06$ T. The solid lines are fits by Eq. (5) (see text).

with the same theory, although none could be measured as accurately as RTS 2.

In the theory, α depends only on the defect potential. For simplicity, let us assume the potential before capture, V, is attractive and radially symmetric, while the potential after capture is zero. The one-electron eigenstates can be classified by the perpendicular component of the angular momentum $m = 0, \pm 1, \pm 2, \ldots$, together with spin and valley indices. In the 2DEG at $B = 0$ there are two equivalent conduction band valleys, and for convenience we combine the spin and valley into a single index $s = 1, \ldots, 4$. If the phase shift for channel (m, s) at E_F in the presence of V is $\delta_{m,s}$, then α is given by [15]

$$\alpha = \left(-1 + \frac{\delta_{0,1}}{\pi}\right)^2 + \sum_{(m,s)\neq(0,1)} \left(\frac{\delta_{m,s}}{\pi}\right)^2. \tag{6}$$

Note that if $V = 0$, $\delta_{m,s} = 0$ for all m and s, $\alpha = 1$, and with the correct value of C Eq. (5) reduces to the noninteracting result, Eq. (3). Charge neutrality forces the phase shifts to obey the Friedel sum rule [16], $\sum_{m,s} \delta_{m,s} = \pi$, and we can use this together with Eq. (6) to obtain a lower limit on α. This limit is reached for pure s-wave scattering, when $\delta_{m,s} = 0$ for all $m \neq 0$. In 2D an attractive potential always has a bound state. When screening is strong, i.e., for large E_F, the bound state may be occupied by four electrons simultaneously [17], and one finds $\alpha \geq 3/4$. For weaker screening (small E_F) the bound state is occupied by only one electron and one finds $\alpha \geq 0$. The limit $\alpha = 0$ is reached when the electron tunnels directly to the defect level from the bound state, while all extended states remain completely unaffected; i.e., all the screening is done by the single electron in the bound state. Then $\delta_{m,s}$ is zero for all channels except the one in which the bound state was destroyed, whose phase shift is $\delta_{0,1} = \pi$.

The value $\alpha = 0.21$ obtained for RTS 2 is only consistent with a singly occupied bound state, with $\pi/2 < \delta_{0,1} < \pi$ and fairly small phase shifts in other channels. On the other hand, the values $\alpha \sim 0.7$ and 0.9 for RTSs 1 and 3, respectively, allow the possibility of multiple occupancy of the bound state. At this point it is interesting to reconsider the results of some earlier RTS experiments in MOSFETs [13,18]. In Ref. [13] the data were fitted using an expression identical in form to Eq. (5) but derived from the theory of two-state systems (TSSs). This remarkable identity is in fact no coincidence, because an electron tunneling between a defect level and a single bound state in the defect potential constitutes a TSS. If one could vary α from 0 to 1 one could, in principle, transform the system continuously from a noninteracting TSS into a noninteracting defect level. While the defects in the present work are in the intermediate regime, those in Ref. [13] were found to have $\alpha \ll 1$, corresponding to TSSs very weakly coupled to the environment (the 2DEG). In this limit Eq. (5) takes the dramatically different form of a narrow Lorentzian centered at $E_d = E_F$. We suggest that for such defects, which are seen only in electrically stressed devices, the 2DEG is locally depleted out due to potential fluctuations at the interface. The extended states therefore remain unaffected on electron capture from the bound state because they are distant from the defect- and bound-state systems. For these defects we can also offer a resolution of a paradox that would arise if the defect were located in a metallic region, namely, that a small value of α implies very weak scattering of electrons at E_F, while a large RTS amplitude appears to require strong scattering. If the defect actually lies in a depleted region then electrons at E_F may be able to tunnel across this region via the bound state when the defect is ionized. Hence the conductance decreases by as much as e^2/h when the defect captures an electron, because the bound state is destroyed, even though all extended-state wave functions are unaltered.

Finally, we examine the effects of a magnetic field on another RTS, RTS 3 (seen in a different device from RTS 1), which lend further support to the basic picture of electrons tunneling between a defect and the 2DEG. For RTS 3, $\alpha \sim 0.9$, so the deviations from the noninteracting result, Eq. (3), are small. In Fig. 3, trace (i) shows the variations of G_1 and G_2 up to 12 T. They are almost identical and were measured separately by sweeping the magnetic field while sampling the conductance at 1 kHz so that the evolution of both levels of the RTS could be seen simultaneously, as illustrated by the sample of such a sweep shown in the inset. The edge of a quantum Hall plateau is visible at the highest field, while only universal conductance fluctuations [19] can be seen at lower B. Trace (ii) shows the corresponding variation of V_{g0}, which above about 4 T undergoes oscillations commensurate with the vertical dotted lines marking points at which n Landau levels are full. This confirms our expectation that the tunneling process should be sensitive to the modulations of the density of states in the 2DEG, although the nature of this sensitivity is complex and will be the subject of future work.

Trace (iii) shows $G_1 - \bar{G}$, where $\bar{G}$ is the smooth monotonic background decrease of G_1 and G_2 visible in trace (i). Trace (iv) shows the RTS amplitude $\delta G_{21} \equiv G_2 - G_1$ over the same field range. As can be seen, the fluctuations in δG_{21} and $G_1 - \bar{G}$ are qualitatively similar and exhibit virtually the same correlation length B_c [19]. In the range $0 < B < 5$ T we find $B_c = 0.04 \pm 0.01$ T, giving an estimate of $A_\phi = (B_c e/h)^{-1} = 0.1\ \mu\text{m}^2$ for the typical phase-coherent area, compared with the device area $A \sim 0.5\ \mu\text{m}^2$. The rms amplitude of the fluctuations is $\sigma_1 = 10 \pm 1\ \mu$S for G_1 and $\sigma_{21} = 3.6 \pm 0.4\ \mu$S for δG_{21}. Hence $\sigma_{21}/\sigma_1 \sim 0.36 \pm 0.06$. Allowing for the uncertainties in A and A_ϕ, this is consistent with the quantum interference prediction $\sigma_{21}/\sigma_1 \sim (A_\phi/A)^{1/2} \sim 0.45$ for the effect of removing a single strong scatterer from a disordered 2D system [2,20]. Trace (iv) also contains

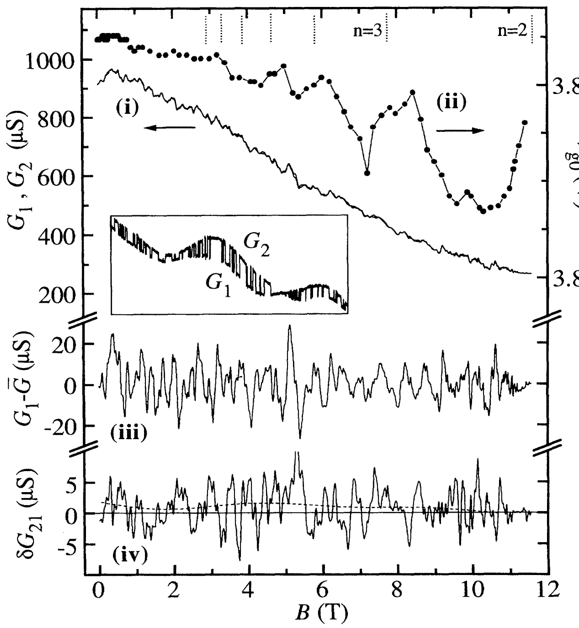

FIG. 3. Effects of magnetic field on RTS 3 at 100 mK, for which $\eta = 0.007$ at $B = 0$. Trace (i) (left axis), conductance G_1 in state 1 of the RTS. Trace (ii) (right axis), balance gate voltage V_{g0}, showing oscillations commensurate with Landau level index n. Inset, sweep of magnetic field over a small range at high bandwidth, where both levels of the RTS are visible. Traces (iii) and (iv), variation of ΔG_1 and δG_{12} over the full field range.

what may be the first evidence for the *classical* effect of removing a scatterer. Classically one would expect the conductance to increase by approximately a fraction $1/N_i$, where N_i is the total number of scatterers. This may be estimated by assuming that the cross section of each scatterer is the Fermi wavelength λ_F, giving $N_i \sim A/\lambda_F l \sim 0.5\ \mu\mathrm{m}^2/(15\ \mathrm{nm} \times 30\ \mathrm{nm}) \sim 1000$. The dashed line superimposed on trace (iv) was obtained by smoothing δG_{21} over a range of 2 T. Notice that it is always above zero: After averaging away the quantum interference fluctuations, the conductance of the occupied state, G_2, is greater than that of the empty state, G_1, by around 1 μS. This is indeed of the order of 1000 times smaller than the total conductance G_1.

In conclusion, using random telegraph signals we have shown that electronic tunneling between a defect level and a 2DEG exhibits the finite-temperature counterpart of the Fermi-edge singularity, with exponent α ranging from nearly zero to nearly 1. Defects with very small values of α behave as two-state systems. We also report magnetic field measurements that support our findings.

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