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Field-effect transistor structures with a quasi–1D channel

S. V. Rotkin†, H. Ruda‡ and *A. Shik*‡

† Beckman Institute for Advanced Science and Technology, University of Illinois at Urbana-Champaign, Urbana, IL 61801, USA

‡ Electronic Material Group, University of Toronto, Toronto M5S3E4, Canada

Abstract. A theory of drift–diffusion transport in a low–dimensional field effect transistor is developed. Two cases of a semiconductor nanowire and a single–wall nanotube are considered using self-consistent electrostatics to obtain a general expression for transconductance. This quantum wire channel device description is shown to differ from classical device theory because of the specific nanowire charge density distribution.

In the present work we consider carrier distributions and parameters of a field effect transistor (FET) with a quantum wire channel, which may be a semiconductor nanowire or a carbon nanotube. The structure includes source and drain electrodes connected by a nanowire/nanotube $-L/2 < x < L/2$, and a gate electrode separated by a thin dielectric layer of the thickness d. We assume the wire to be uniformly doped with 1D density $N = \mathrm{const}(x)$. When the structure is in operation, the source-drain voltage V_d causes a current j along the channel and a re-distribution of carrier concentration as compared with the initial specific density. A voltage V_g is applied to the gate and changes the concentration, which controls the FET transport.

All potentials are measured from the middle point of the wire ($x = 0$) so that the source and drain potentials are $-V_d/2$ and $V_d/2$. In this case the potentials along the wire and concentration changes caused by V_g together with the contact potentials, and by V_d are, respectively, symmetric and antisymmetric functions of x and are given the subscripts s and a: $\phi_{s,a}(x)$ and $n_{s,a}(x)$.

The potentials $\phi_{s,a}(x)$ can be divided into two parts: the components $\phi^0_{s,a}(x)$ created by electrodes and contact potentials, which should be found from the Laplace equation containing no channel charge density and the components $\phi^1_{s,a}(x)$ caused by the electron charge in the channel $-en_{s,a}(x)$. We assume that the characteristic lengths L and d determining the potential and density distribution along the channel, noticeably exceed the nanowire/nanotube radius a. In this case the relationship between $\phi^1_{s,a}(x)$ and $n_{s,a}(x)$ is approximately linear [1–4] and for a nanowire with non-degenerate carriers, the current j can be written as:

$$\frac{j}{e\mu} = n(x)\frac{d\phi^0}{dx} - \left[\frac{2e}{\varepsilon}\ln\left(\frac{l}{a}\right)n(x) + \frac{kT}{e}\right]\frac{dn}{dx} \qquad (1)$$

where $n = n_s + n_a$, $\phi^0 = \phi^0_s + \phi^0_a$, ε is the ambient dielectric susceptibility, μ is the carrier mobility and l has the order of $\min\{L, 2d\}$.The second term describes the drift in the self-consistent field $-\nabla\phi^1_{s,a}(x)$ and the last term corresponds to the diffusion current, which in one–dimensional transport cannot be neglected in comparison with the drift component. For a nanotube with degenerate carriers, the thermal energy kT should be replaced by the Fermi energy and (1) reads as:

$$\frac{j}{e\mu} = n(x)\frac{d\phi^0}{dx} - eC_t^{-1}n(x)\frac{dn}{dx}, \qquad (2)$$

where $C_t^{-1} = C_g^{-1} + C_Q^{-1}$ is the inverse capacitance of the nanotube derived in [3] and containing both logarithmic geometrical capacitance, $C_g^{-1} \sim \ln(l/a)$, similar to (1) and the quantum capacitance of the 1D electron gas, $C_Q = 1/(e^2\nu) \simeq 0.31$ for one degenerate subband of a single wall nanotube. The same simplified expression is also derived from (1) in the limit $A = 2e^2N\ln(l/a)/(\varepsilon kT) \gg 1$.

The boundary conditions $n(\pm L/2) = n_c$ assume the source and drain to support constant concentration at the contacts, independent of the applied voltage. For $n_c \neq N$ our formulae require some modification. In the closest vicinity of contacts a finite charge density $e(N - n_c)$ exists in the channel. To provide equipotentiality of metallic contacts, we must assume the presence of oppositely charged images inside the contacts. But this means discontinuity of charge density at $x = \pm L/2$ and makes doubtful the adequacy of continuum approach assuming smooth charge and potential variations. To avoid this difficulty, we measure n from n_c by assuming $n(x) = n_c + \Delta n(x)$ with the simultaneous inclusion in potential the component $\phi_c(x)$ representing potential of a wire with uniform charge $e(N - n_c)$ between metallic contacts $x = \pm L/2$ calculated in [2].

Now we can find the carrier concentration $n(x)$ and the electric current j caused by the voltages V_g and V_d. The problem is relatively simple if we restrict ourselves to the linear case of small V_d. In the zeroth approximation both $\phi^0(x)$ and $n(x)$ contain only a symmetric component: $n_s(x) = n_c + \Delta n_s(x)$. $\Delta n_s(x)$ should be found from (1) with $j = 0$ and $\phi^0(x) = \phi_c(x) + \phi_g(x)$ where $\phi_g(x)$ is the potential created by the gate electrode. When $n_s(x)$ is found, (1) can be linearized in n_a and solved with the conditions $n_a(0) = n_a(L/2) = 0$. This gives the concentration profile and the current j, which appears especially simple for a single-wall nanotube or $A \gg 1$ where the problem is reduced to the ordinary Kirchhoff's law:

$$j = \frac{V_d}{R}, \qquad R = \frac{2}{e\mu}\int_0^{L/2}\frac{dx}{[n_c + \Delta n_s(x)]}. \qquad (3)$$

The resulting j depends on the gate voltage V_g through the functions $\phi^0(x)$ and $\Delta n_s(x)$, which allows us to calculate the transistor transconductance dj/dV_g.

The potential profile $\phi^0(x)$ and hence all the results depend noticeably on the geometry of the source and drain contacts. We consider in detail bulk contacts representing metallic or heavily doped semiconductor regions with all three dimensions considerably exceeding a, d and L and thus assumed to be infinite. By solution of corresponding Laplace equation we find the potential distribution $\Phi(x, y)$ created by this system

of electrodes in the absence of a wire. The resulting $\Phi(x,d)$ contains two terms: that proportional to V_g and corresponding to $\phi_g(x)$ and that proportional to V_d which corresponds to $\phi_a^0(x)$.

If the simple expression (3) is applicable, the dimensionless current $i = jL/(n_0 e\mu V_d)$ has an explicit form:

$$i = \left[2\int_0^{1/2} \frac{dt}{1+g\Psi(t)}\right]^{-1} \qquad (4)$$

where $g = 2\varepsilon V_g/[\pi e n_0 \ln(l/a)]$ is the dimensionless gate voltage and for $n_c = n_0$,

$$\Psi(t) = \sum_0^\infty \frac{(-1)^n}{(2n+1)} \cos[\pi t(2n+1)] \exp[-\pi d(2n+1)/L].$$

The channel conductivity near the cut-off is determined by the point of minimal concentration, which in a symmetric structure is $x = 0$, and hence is determined by the properties of $\Psi(t)$ at small t. Expansion of $\Psi(t)$ allows us to perform integration in (4) and obtain

$$\begin{aligned} i &= \frac{\sqrt{(g-g_0)\sinh(\pi d/L)}}{\sqrt{2}\cosh(\pi d/L)}, \\ g_0 &= -\left\{\frac{\pi}{2} - \arctan[\exp(-\pi d/L)]\right\}^{-1}. \end{aligned} \qquad (5)$$

Thus the transconductance di/dg increases in the vicinity of cut-off $\sim (g-g_0)^{-1/2}$.

In the structures with two- and even quasi-one-dimensional contacts the potential profiles ϕ_c, ϕ_g and ϕ_a have essentially different x-dependence with singularities near contacts [2]. In our case it means a different profile for the $\Psi(t)$ function. Its expansion near the maximum is, of course, also quadratic and hence results in the same qualitative final result $di/dg = B(g-g_0)^{-1/2}$, with particular values of the cut-off voltage g_0 and the coefficient B different from the case of bulk contacts.

The simplified expressions (2), (4) neglected diffusion, which is equivalent to the limit of zero temperature. The formula $\Delta n_0(x) = C_t\phi_c(x)$ resulted from (2) in equilibrium simply gives $n = 0$ for all points where $\phi^0(x) < -C_t^{-1}n_c$. Thus in the linear approximation, the cut-off voltage g_0 corresponds to the condition $\phi^0(0) = -C_t^{-1}n_c$ and at lower g the current is exactly zero. It is evident that at non-zero temperatures the current at $g < g_0$ will have an activation behaviour: $j \sim \exp(-\Delta/kT)$ where $\Delta = e(-C_t^{-1}n_c - \phi^0(0))$. Since $\phi^0(0)$ depends linearly on V_g, the activation energy Δ is directly proportional to $g_0 - g$.

In the case of arbitrary V_d when the linear approach fails, the problem requires numerical solution of the general non-linear equations (1) or (2) with the potential consisting of three parts: $\phi^0(x) = \phi_c(x) + \phi_g(x) + \phi_a(x)$ describing, respectively, the influence of contact work function, gate voltage and source-drain voltage and calculated earlier. Two boundary conditions: $n(\pm L/2) = n_c$ determine the integration constant and the so far unknown value of j. Since $\phi_g(x)$ is proportional to V_d and $\phi_a(x)$ is proportional to V_d, the resulting solution gives us the current-voltage characteristic (CVC) $j(V_d)$ for various gate voltages.

To calculate it quantitatively, we choose particular values $\ln(l/a) = 3$ and $d/L = 0.3$. As an illustration, we consider the ideal Ohmic contacts with $n_c = N$ where the dimensionless threshold estimated with (5) $g_0 = -2.7$. Figure 1(a) shows CVC at $g = -2.8$ (below the threshold) and $g = -2.5$ (above the threshold). All characteristics have a superlinear character, which has a simple explanation. High drive voltage V_d tends to distribute carriers uniformly along the channel. In our conditions when powerful contact reservoirs fix the concentration n at the points where it is maximal, such a re-distribution will increase the minimal value n in the center of channel and hence increase conductivity of the latter.

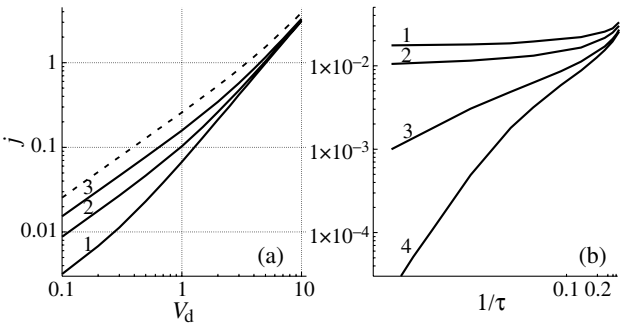

Fig. 1. (a) CVC of a nanowire with $d/L = 0.3$ at $V_g = -2.5$ (dashed line) and $V_g = -2.8$ (solid lines) for $\tau = 0.05$ (1), 0.1 (2), and 0.2 (3). (b) temperature dependence of the current at $V_g = -2.5$ (1), -2.6 (2), -2.7 (3), -2.8 (4).

Figure 1(a) presents also information on temperature dependence of the channel conductivity. Above the threshold this dependence is practically absent. The CVC curves for $g = -2.5$ at different temperatures does not deviate from the dashed line corresponding to $\varepsilon kT/(e^2 n_c) = 0.2$ more than by 10% and are not shown in the figure. For V_g below threshold and not very high V_d, Fig. 1(a) demonstrates a strong temperature dependence of the current shown in more details in Fig. 1(b) calculated for low enough V_d corresponding to the linear initial part of CVC. While the two upper curves for the above-threshold V_g have no noticeable temperature dependence, the two lower curves demonstrate such a dependence with Δ growing with $|V_g|$, in accordance with the predictions. At high V_d, where contact injection and electric field tend to create uniform carrier concentration equal to n_c, the different curves merge and temperature dependence collapses.

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