Direct Solution of the Boltzmann Transport Equation and Poisson–Schrödinger Equation for Nanoscale MOSFET

Stefano Scaldaferrri
Dialog Semiconductor GmbH, Kirchheim unter Teck-Nabern

Gilberto Curatola
NXP Semiconductors, Leuven

Giuseppe Iannaccone
Dipartimento di Ingegneria dell’Informazione: Elettronica, Informatica, Telecomunicazioni, Università di Pisa

Abstract—We propose an efficient and fast algorithm to solve the coupled Poisson–Schrödinger and Boltzmann transport equations (BTE) in two dimensions. The BTE is solved in the relaxation time approximation within each subband obtained from the direct solution of the Schrödinger equation. The proposed approach, considering a subband-based transport formalism, allows to fully explore the entire range from drift-diffusion to ballistic regime in nanoscale field-effect transistors. Quantum effects are also fully taken into account by the direct solution of the Schrödinger equation. The model is implemented in the NanoTCAD2D device simulator and used to study the device performance of a 25-nm channel-length MOSFET. The influence of scattering on the electron distribution function and on device characteristics is analyzed in detail.

Index Terms—Ballistic transport, Boltzmann transport equation, modeling, nanoscale MOSFET device, Schrödinger equation.

I. INTRODUCTION

THE CONTINUOUS scaling of device dimensions to meet the requirements in terms of silicon area occupation and device performance that is set by the Semiconductor Industry has led to a rapid increase of the complexity of CMOS technology and, consequently, of the TCAD tools needed to obtain reliable simulations.

Monte Carlo techniques for the solution of the Boltzmann transport equation (BTE) offer a high degree of accuracy and allow to properly describe the essential physics of nanoscale devices even in strong off-equilibrium conditions [1]–[3]. Unfortunately, the huge computational resources required are among the reasons why the vast majority of routinely performed simulations are still based on simplified approaches, like the drift-diffusion (DD) or hydrodynamic (HD)-type models [4], [5].

However, all the simplified models (DD and HD) rely on strong approximations on the carrier distribution function, which is necessary to make the problem analytically manageable. Such approximations may invalidate the results of simulations in the case of ultrascaled devices and/or strong off-equilibrium conditions; moreover, the BTE [6]–[8] is strictly semiclassical and does not allow to take into account quantum effects which become predominant as the dimensions of the device enter the sub-100-nm regime.

Even if, recently, some Monte Carlo tools [9]–[11] have been modified in order to include the quantum effects by means of approximate models [15]–[18] or by direct coupling the Schrödinger equation to the Boltzmann equation, the computing time is still excessive for industrial application, and therefore, such approaches are confined to the computational electronics community.

We will show in the following that in the case of confined systems, instead, where the dimensionality is reduced due to the quantum confinement, it is possible to directly solve the BTE (i.e., without a statistical approach) and to couple it to the Schrödinger equation without large penalty in the computational efficiency of the algorithm.

Previous attempts in this direction have been presented in the literature by Goldsman et al. [12]–[14]. We want to stress here that our main goal is to develop a 2-D simulator for nanoscale MOSFETs that satisfies the following requirements.

1) It self-consistently solves the coupled Poisson–Schrödinger–Boltzmann equation in two dimensions without a priori assumptions on the carrier distribution functions.
2) Computing time is maintained comparable to the commercial DD or HD tools.
3) It accounts for mass anisotropy (the Schrödinger equation is solved three times for the electrons and twice for holes).
4) It can be applied to a wide variety of semiconductor materials, like Si, Ge, AlGaAs, SiGe, and strained-Si.
5) It includes a fully ballistic model for electrons and holes.
6) Differently from other “approximate” quantum models proposed in the literature, our approach allows to have full information on the subband structure. Therefore, it permits a detailed evaluation of the different scattering rates in each subband and among different subbands given the knowledge of the envelope wavefunctions at each position along the channel.
As far as scattering is concerned, we present here the results of simulations performed in the simplified assumption of negligible intersubband scattering which gives us the freedom of treating electron transport in each subband independently. Additional results obtained with a more detailed scattering model will be presented elsewhere [19]. Here, we are mainly interested in demonstrating the feasibility of the proposed approach.

This paper is organized as follows. In Section II-A, we present the algorithm for the solution of the BTE, assuming a known subband profile and in the relaxation time approximation (RTA). For simplicity, we also assume a scattering time independent of energy, but the algorithm can be easily generalized. The full self-consistent Poisson–Schrödinger–Boltzmann approach is then described in detail in Section II-B.

In Section III, we highlight the numerical issues connected with the discretized form of the BTE.

Finally, in Section IV, we present the results of simulations performed on a benchmark nanoscale transistor, i.e., the “well-tempered” MOSFET with an effective channel length of 25 nm proposed by Antoniadis et al. [20]. Electrical characteristics are presented for a whole range of operating regimes ranging from DD to fully ballistic. Furthermore, details on the electron distribution function in each subband only depends on the $y$ coordinate and of the wave vector along directions $y$ and $z$.

Under these assumptions, the steady-state Boltzmann equation can be written as

\[ \nu_y \frac{\partial f(y, k)}{\partial y} - \frac{q \xi_y}{\hbar} \frac{\partial f(y, k)}{\partial k_y} = \Theta(f) \]  

(1)

where $k = k_y \hat{y} + k_z \hat{z}$ represents the electron wave vector in the transport plane, $q$ is the electron charge, $\hbar$ is the reduced Planck constant, $\nu_y$ is the $y$-component of the electrons’ group velocity, $\xi_y$ is the $y$-component of electric field, and $f$ is the electron distribution function in the phase space of the considered subband.

The electric field in the transport direction can easily be evaluated given the knowledge of the subband profile and reads

\[ \xi_y = \frac{1}{q} \frac{dE_{SB}(y)}{dy} . \]  

(2)

The operator $\Theta(f)$ at the right-hand side of (1) represents the collision term, which can be expressed as

\[ \Theta(f) = \sum_{k'} S(k'; k) f(y, k') \left[ 1 - f(y, k) \right] + \sum_{k'} S(k; k') f(y, k) \left[ 1 - f(y, k') \right] \]  

(3)

where $S(k'; k)$ represents the scattering rate from state $k'$ to state $k$. 

II. Model

Let us consider the nMOSFET shown in Fig. 1, where $x$ represents the direction of quantum confinement, and $y$ represents the direction of electron transport. Electrons moving from source to drain must overcome an energy barrier whose height is modulated via the voltage applied to the gate electrode. As a consequence of gate-induced vertical field, carriers are confined at the oxide/silicon interface in condition of strong inversion. This causes energy quantization in the transverse direction.

The position (in energy) of the available states obtained from the solution of the Schrödinger equation varies along the current flow direction due to the varying confining potential along the channel. Available states are therefore grouped in different subbands. A generic subband profile along the current flow direction is shown in Fig. 2, where $E_{SB}(y)$ represents the subband edge at coordinate $y$.

The Fermi levels at the two boundaries are set by the source and drain contacts $E_{FS}$ and $E_{FD}$, respectively. Indeed, we assume that left- and right-moving electrons are injected in the channel with the chemical potential of the emitting reservoir.

Considering the steady-state Boltzmann equation, for each subband, we assume that the kinetic term along $x$ is imposed by energy quantization, and therefore, the distribution function in each subband only depends on the $y$ coordinate and of the wave vector along directions $y$ and $z$. 

Under these assumptions, the steady-state Boltzmann equation can be written as
A. One-Subband Transport

Under the RTA hypothesis, the scattering term can be simplified in the form
\[
\Theta(f) = \frac{f_{\text{eq}}[E_{\text{in}}(y), k] - f(y, k)}{\tau(k)}.
\]

The term \( f_{\text{eq}}[E_{\text{in}}(y), k] \) represents an equilibrium distribution function for the 2-D electron gas, i.e., a Fermi–Dirac distribution. \( E_{\text{in}}(y) \) represents the quasi-Fermi level at coordinate \( y \), and \( \tau(k) \) is the relaxation time which represents the time required for the carrier population to reach an equilibrium distribution (thus, described by a Fermi–Dirac distribution function) as a consequence of the scattering processes.

The equilibrium distribution \( f_{\text{eq}} \) can conventionally be expressed as
\[
f_{\text{eq}}[E_{\text{in}}(y), k] = \left[ 1 + \exp\left(\frac{E_{\text{SB}}(y) + \frac{\hbar^2 k_y^2}{2m_y} + \frac{\hbar^2 k_z^2}{2m_z} - E_{\text{in}}(y)}{k_B T}\right)\right]^{-1}
\]  

where \( m_y \) and \( m_z \) are the electron effective masses in \( y \)- and \( z \)-directions, \( k_B \) is the Boltzmann constant, and \( T \) is the temperature.

The RTA approximation must, however, ensure charge conservation, as expressed by the following relation (in the present form, the following relation strictly holds when the intersubband scattering is neglected):
\[
\int_{k} \frac{1}{2\pi^2} f(y, k) dk = \int_{k} \frac{1}{2\pi^2} f_{\text{eq}}(E_{\text{in}}(y), k) dk.
\]

Previous simulations performed on the same device have shown indeed that in strong inversion conditions, almost 95% of the total carrier concentration resides in the first subband [21]. This result justifies the approximation of negligible intersubband scattering and gives us strong confidence in the simulation results.

The left side of (6) represents the effective carrier concentration \( n(y) \) associated to the considered subband, while the right side represents the equivalent concentration at equilibrium which is a known nonlinear function of the electron quasi-Fermi level and is given by
\[
n_{\text{eq}}(y) = g_e k_B T \sqrt{\frac{m_y m_z}{\pi \hbar^2}} \ln \left[ 1 + \exp\left(\frac{E_{\text{in}}(y) - E_{\text{SB}}(y)}{k_B T}\right)\right]
\]  

where \( g_e \) represents the degeneracy of each group of equivalent conduction band minima of silicon (in our case, \( g_e = 2 \)).

In particular, the charge-conservation property gives us an important relation between the equilibrium and quasi-Fermi levels out of equilibrium. Inverting (6) with the use of (7) allows us to express the quasi-Fermi level \( E_{\text{in}}(y) \) in the considered subband as
\[
E_{\text{in}}(y) = g(n(y))
\]  

where \( g \) is a known nonlinear function.

In the rest of this paper, we will assume for simplicity a relaxation time constant with energy. The extension to an energy-dependent relaxation time is, however, straightforward. Under this hypothesis, (1) can be rewritten as follows:
\[
v_y \frac{\partial f(y, k)}{\partial y} - \frac{q \xi_y}{\hbar} \frac{\partial f(y, k)}{\partial k_y} = \frac{f_{\text{eq}}(E_{\text{in}}(y), k) - f(y, k)}{\tau(y, k)}
\]

If we define
\[
F(y, k_y) \equiv \int_{k_x} \frac{1}{2\pi^2} f(y, k) dk_x
\]

we can integrate both terms of (9) in \( k_z \) and obtain
\[
v_y \frac{\partial F(y, k_y)}{\partial y} - \frac{q \xi_y}{\hbar} \frac{\partial F(y, k_y)}{\partial k_y} = \frac{F_{\text{eq}}(E_{\text{in}}(y), k_y) - F(y, k_y)}{\tau(y, k_y)}.
\]

The term \( F(y, k_y) \) represents the derivative of the electron distribution function with respect to the \( y \)-component of the wave vector. Indeed, by definition
\[
\int F(y, k_y) dk_y = n(y).
\]

The computation of \( F_{\text{eq}}(E_{\text{in}}(y), k_y) \) from the quasi-Fermi levels requires instead the evaluation of Fermi integral on the order of \((-1/2)\)
\[
F_{\text{eq}}(E_{\text{in}}(y), k_y) = \sqrt{\frac{2m_z k_B T}{\pi \hbar^2}} F_{-\frac{1}{2}} \left(\Delta E \right)
\]

where
\[
\Delta E \equiv E_{\text{in}}(y) - \left( E_{\text{SB}}(y) + \frac{\hbar^2 k_y^2}{2m_y} \right)
\]

\[
F_{-\frac{1}{2}}(x) = \int_{y=0}^{\infty} \frac{y^{-\frac{1}{2}}}{1 + \exp(y - x)} dy.
\]

With the previous assumptions, the original problem has been reduced to a nonlinear system composed by (11) and (8) whose unknowns are the distribution functions \( F(y, k_y) \) and the quasi-Fermi levels \( E_{\text{in}}(y) \). Furthermore, the complexity of the problem has been reduced with respect to the original case since the discretization domain is now a 2-D space whose components are the electron transport direction \( y \), in real space, and the \( y \)-component of the wave vector, \( k_y \), in \( k \) space.

A box integration method [22] is used to discretize (11) and (8) onto nonuniform rectangular grids. If we indicate the number of internal grid points with \( N \), the result of discretization is a nonlinear algebraic system
\[
A \cdot \mathbf{F} = \frac{\mathbf{F}_{\text{eq}}}{\tau}
\]
At each step, the potential profile obtained from the PS solver is used to evaluate the new guess of the distribution function which is then used to estimate the carrier density. This term is inserted afterward into the Poisson equation, and the self-consistent procedure continues until convergence is reached.

III. NUMERICAL ISSUES

In this section, we will focus on the numerical issues arising from the discretization of the BTE. First, we specify boundary conditions on the integration domain, and then, we identify an appropriate discretization scheme.

A. Boundary Conditions

The integration domain of (11) is the $y-k_y$ plane. In the real-space dimension, the boundary conditions are set by source and drain contacts, while in the momentum space, we can properly fix an interval $[k_{y\text{min}} \div k_{y\text{max}}]$ outside of which the electron distribution function can be considered negligible. Hence, we can impose Dirichlet boundary conditions on the momentum space assuming the distribution function to be zero for $k_y = k_{y\text{max}}$ and $k_y = k_{y\text{min}}$.

The imposition of boundary conditions in the real space requires instead more subtle argumentations: BTE has indeed only the first derivative in space, but we would specify boundary conditions at both contacts. This seeming inconsistency can be overcome assuming the hypothesis of “black body” contacts. In particular, distribution-function values are specified at each contact only for emitted electrons. An equilibrium distribution function with Fermi level $E_{FS}$ ($E_{FD}$) is therefore fixed at the source (drain) electrode for states characterized by $k_y > 0$ ($k_y < 0$).

B. Discretization

In order to discretize (11), both dimensions of the integration area can be divided into nonuniform steps. If we indicate the number of points along $y$ and $k_y$ axes with $N_y$ and $N_{k_y}$, we can identify the coordinates of real space and momentum-space dimensions with $y[i]$ and $k_y[j]$, respectively, where $i = 1, \ldots, N_y$, and $j = 1, \ldots, N_{k_y}$.

Let $(i,j)$ be the grid point of coordinates $(y_i, k_{yj})$ and $F_{i,j}$ be the electron distribution function at that point. According to a box integration approach, the device domain is divided into as many nonoverlapping cells as internal grid points.

Integrating (11) on each cell leads to the following finite difference equation:

$$v_{yj} [F_{i+\frac{1}{2},j} - F_{i-\frac{1}{2},j}] \Delta k_{yj} - \frac{q\xi_{yj}}{\hbar} [F_{i,j+\frac{1}{2}} - F_{i,j-\frac{1}{2}}] \Delta y_i$$

$$= \frac{F_{eq,i,j} - F_{i,j}}{\tau} \Delta k_{yj} \Delta y_i \quad (16)$$

where subscripts that are applied to physical quantities denote coordinates of the domain space, while $\Delta k_{yj}$ and $\Delta y_i$ represent the dimensions of the elementary cell associated with the $(i,j)$ grid point. Equation (16) is a balance of fluxes in the control
volume: The first two terms at the left-hand side are diffusion and drift fluxes, respectively, while the right-hand side is the scattering flux. An upwinding technique is used to obtain a stable discretization and to avoid the unphysical negative values for $F$.

IV. SIMULATION RESULTS

The method described in Section II has been used to study the electrical characteristics of a benchmark device, i.e., the “well-tempered” MOSFET proposed by Antoniadis et al. [20]. In spite of the reduced dimensions, short channel effects (SCEs) in the device are maintained at a reasonable level, owing to the “ideal” doping profile considered. In particular, Fig. 5 shows the net doping concentration in the transistor: Source and drain profiles are Gaussian, while a superhalo doping is implanted in the substrate in order to reduce SCEs. The oxide thickness is 1.5 nm, and the effective channel length, which is measured at the points where minority carrier concentration in the channel falls below $2 \times 10^{19}$ cm$^{-3}$, is approximately 25 nm. The polysilicon gate has a donor concentration of $5 \times 10^{20}$ cm$^{-3}$ to reduce polysilicon depletion effects.

Transport regime of carriers in the channel can be selected by specifying the scattering time in the input of the BTE solver. In the present simulations, we choose a value in the range $[10^{-14}$ s $\div 10^{-11}$ s], which allows us to explore the entire DD to ballistic transport regime.

For the BTE solution, we have adopted a grid with $N_y = 951$ points in the real-space dimension and $N_k = 1000$ points in the momentum dimension.

In order to improve the stability of the algorithm described in Section II, an underrelaxation technique has been used to update the subband profiles at each different iteration step. By indicating the subband profiles at two successive iteration steps obtained from the PS solver with $E_{k-i}^i$ and $E_{k-i}^{i-1}$, the new subband profile $E_{k-i}^{\text{new}}$ for the BTE solver is obtained as

$$E_{k-i}^{\text{new}} = E_{k-i}^{i-1} + \beta (E_{k-i}^{k-1} - E_{k-i}^{i-1}).$$ (17)

The stability of the implemented algorithm and its convergence properties is first analyzed in Figs. 6 and 7. In particular, Fig. 6 shows the residual versus the number of iterations for a bias condition ($V_{gs} = 1$ V, $V_{ds} = 1$ V, and $\tau = 10^{-13}$ s).

As is clearly visible, the convergence rate strongly depends on the bias conditions applied to the structure. Fig. 7 shows the iteration number as a function of the voltage applied to the gate electrode $V_{gs}$ for $V_{ds} = 1$ V and a scattering time $\tau = 10^{-13}$ s. In accordance with the residual behavior, the time required to reach convergence increases with an increasing bias.

The effect of carrier scattering on the electrical characteristics of the device is analyzed in the following for a fixed drain–source voltage $V_{ds} = 1$ V and a gate bias sweeping from $V_{gs} = 0.3$ V to $V_{gs} = 1.3$ V. Fig. 8 shows the transfer
characteristics in semilogarithmic scale for scattering times $\tau = 10^{-11}$ s, $10^{-12}$ s, $10^{-13}$ s, and $10^{-14}$ s.

The threshold voltage $V_T$ for the considered structure is about 0.65 V. This high value is the consequence of the very high halo doping concentration required to suppress the SCEs arising from the protruding of the source and drain depletion regions into the channel. It is worth noting that in a real device, the \textit{ad hoc} doping profile considered would induce a severe current-drive capability reduction as well as strong band-to-band tunneling effects at the junctions.

Subthreshold slope is close to 100 mV/decade, and as expected, it is barely affected by the scattering time $\tau$ selected. Conversely, the current drive is strongly dependent on the scattering conditions. As shown in Fig. 8, indeed, for a high scattering time set as input ($\tau = 10^{-12}$ s), the current approaches the theoretical ballistic limit, while a strong current degradation can be observed as the scattering time is progressively decreased down to $\tau = 10^{-14}$ s.

Fig. 9 shows the distribution-function color maps related to the first subband for high bias condition ($V_{gs} = 1.3$ V, and $V_{ds} = 1$ V). The four cases analyzed roughly correspond to a transition from a full DD regime ($\tau = 10^{-14}$ s) to a full ballistic case ($\tau = 10^{-11}$ s).

In the BB case, as theoretically predicted [26]–[28], we can clearly observe the ballistic peak in the drain region and a hemi-Fermi–Dirac distribution-function profile at the top of the barrier due to the suppression of electrons coming from the drain contact. Indeed, in the case of high source-to-drain bias, carriers injected from the drain reservoir do not have enough energy to overcome the energy barrier set by $V_{DS}$ and cannot reach the source contact. The distribution function at the source contact is therefore represented by a hemi-Fermi–Dirac distribution.

At the other extreme, going toward a DD regime ($\tau = 10^{-14}$ s), collision randomize carriers motion, and the distribution function has a quasi-equilibrium behavior at each channel position.

It is worth noting that while the terminal current is not strongly dependent on $\tau$ in the range $[10^{-11}−10^{-13}$ s], the carrier distribution shows instead a strong dependence on the scattering rate, particularly in the polysource implant superposition region. Conversely, for $\tau = 10^{-14}$ s, a change in the carrier distribution directly translates into a reduced terminal current.

This behavior can be explained with the help of Fig. 10, where the first subband profile along the channel is shown for each scattering time considered. The top of the barrier is located in the source-contact proximity, while accumulation regions for $y$ in the range $[-30 \text{ nm} \div -15 \text{ nm}]$ and $[15 \text{ nm} \div 30 \text{ nm}]$ are the effects of the superposition of polygate with source and drain extensions, respectively.

The maximum of the barrier is practically coincident with the source contact. For $\tau > 10^{-13}$ s, almost all source emitted electrons cross the barrier, and the barrier peak is barely affected by a change in $\tau$. Any scattering event occurring after the maximum point is then revealed by the presence of a quasi-symmetric electron fraction but does not affect the current at the terminals.

For $\tau = 10^{-14}$ s, instead, scattering directly influences the barrier height which modifies the electron fraction overcoming the barrier and, consequently, the terminal current.

A. Ballistic Fraction

As observed in the previous section, a ballistic transport regime can be identified by a peak in the electron distribution functions in the drain region which gradually vanishes when approaching a collision-dominated DD regime. This property can be used to define a synthetic parameter indicating the transport regime for a given structure.

Let us refer to Fig. 11 where the electron distribution function $F(y_d, v_y)$ at the drain contact ($y_d = 47.5$ nm) is represented for bias conditions ($V_{ds} = 1$ V, and $V_{gs} = 1.3$ V) and a scattering time ($\tau = 10^{-13}$ s).

We can identify a quasi-symmetric component of the distribution function $F_s(y_d, v_y)$ defined as follows:

$$F_s(y_d, v_y) = \begin{cases} F(y, v_y), & \text{for } -\infty < v_y < v_{y\text{min}} \\ 0, & \text{otherwise} \end{cases}$$

(18) and a ballistic component $F_b(y_d, v_y)$

$$F_b(y_d, v_y) = \begin{cases} 0, & \text{for } -\infty < v_y < v_{y\text{min}} \\ F(y, v_y), & \text{otherwise} \end{cases}$$

(19)
Fig. 9. Electron distribution function in the first subband as a function of four different scattering times for $V_{gs} = 1.3$ V and $V_{ds} = 1$ V. For an operating regime ranging from fully (top left) ballistic to (bottom right) DD, it is clearly visible that the ballistic peak at the drain side progressively disappears due to the randomizing effect of collisions.

Fig. 10. Energy profile of the first subband as a function of position along the channel and the scattering time. In particular, different scattering times have been considered: $\tau = 10^{-11}$ s, $10^{-12}$ s, $10^{-13}$ s, and $10^{-14}$ s.

Fig. 11. Electron distribution function at the drain contact for $V_{ds} = 1$ V, $V_{gs} = 1.3$ V, and $\tau = 10^{-13}$ s. The ballistic and drift-diffusive components are highlighted.

where $v_{y_{\text{min}}}$, as shown in Fig. 11, ideally represents the separation point between the ballistic and quasi-symmetric components. From (18) and (19), the following relation naturally follows:

$$F(y_d, v_y) = F_0(y_d, v_y) + F_s(y_d, v_y). \quad (20)$$

In a similar way, we can identify a drift-diffusive component of the current at the terminals $J_d$ defined as follows:

$$J_d = \int_{k_y} F_s(y_d, v_y) dv_y. \quad (21)$$

If we indicate the total current at the terminals with $J_T$, then the term $B_d = 1 - J_d / J_T$ can represent a synthetic parameter indicating transport regime, which we will call ballistic fraction ($1 - B_d$ represents instead the diffusive fraction).

Fig. 12 shows the diffusive and ballistic fractions $r$ plotted versus the scattering time for $V_{ds} = 1$ V, and $V_{gs} = 1.3$ V.

As we expected, $R_d$ reaches a saturation point close to unity for low scattering-time values (drift-diffusive regime), while it exhibits a linear decreasing profile for the increasing values of scattering time.
V. CONCLUSION

We have presented an algorithm for the self-consistent solution of the BTE and the PS equation in two dimensions, which accurately captures the essential physics of ultrascaled field-effect transistors. The proposed approach, based on a subband-transport formalism, allows to account for the quantum effects due to the gate-induced vertical confining field. Furthermore, the direct solution of the BTE in each subband that is obtained from the direct solution of the Schrödinger equation coupled to a PS solver allows to fully explore the entire DD to ballistic (BB) regime, which is a key capability for modern device simulators.

The proposed method has been then used to study a benchmark device [20] with a 25-nm channel length. In particular, the carrier distribution function and the terminal currents have been carefully analyzed as a function of the scattering events in the device. The main features of the DD and BB transports have been demonstrated as well as the capability of modeling the intermediate operating regime. Finally, a synthetic parameter, indicating the transport regime and the ability of scattering events to alter the terminal current, has been defined and evaluated as a function of the relaxation time.

REFERENCES


Stefano Scaldaferrri was born in Lagonegro, Italy, on November 30, 1978. He received the M.Sc. degree in electrical engineering from the University of Pisa, Pisa, Italy, in 2003, with a thesis on modeling of electronic transport in nanoscale field-effect transistors, and the Master of Engineering degree in embedded systems design from the ALaRI Institute of Lugano, Lugano, Switzerland, in 2004, with a thesis on application-driven optimization of VLIW architectures.

From September 2004 to September 2006, he was with the Infineon Technologies, Munich, Germany, working as an IC Design Engineer for memory products. Since September 2007, he has been with the Dialog Semiconductor GmbH, Kirchheim unter Teck-Naber, Germany, as an IC Designer for audio and power management applications.
Gilberto Curatola was born in Italy, on May 5, 1975. He received the M.Sc. degree in electrical engineering and the Ph.D. degree for a thesis on the subject of quantum effects and transport in nanoscale field-effect transistors from the University of Pisa, Pisa, Italy, in 2000 and 2005, respectively.

Since January 2005, he has been with the NXP Semiconductors, Leuven, Belgium, working on device simulations for CMOS applications. His research interest includes characterization and simulation of deep-submicrometer silicon devices with emphasis on quantum effects and mobility.

Giuseppe Iannaccone received the M.S. and Ph.D. degrees in electrical engineering from the Università di Pisa, Pisa, Italy, in 1992 and 1996, respectively.

In 1996, he took a permanent position as a Researcher with the Italian National Research Council, and in the same year, he obtained a faculty position with the Dipartimento di Ingegneria dell’Informazione: Elettronica, Informatica, Telecomunicazioni, University of Pisa, first as an Assistant Professor and, then, since January 2001, as an Associate Professor of electronics. He has authored or coauthored more than 110 papers published in peer-reviewed journals and more than 70 papers in proceedings of international conferences. His research interests include transport and noise in nanoelectronic and mesoscopic devices, development of device modeling and TCAD tools, and the design of extremely low-power circuits and systems for radio-frequency identification and ambient intelligence scenarios.

Dr. Iannaccone has coordinated a few European and National Projects involving multiple partners and has acted as Principal Investigator in several research projects funded by public agencies at the European and National levels, and by private organizations. He is in the technical committee of few international conferences and serves as a referee for the leading journals in the fields of condensed matter physics, device electronics, and circuit design.