NUMERICAL SIMULATION OF MOS TRANSISTORS

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Abstract. This contribution is intended to review the international state-of-the-art in numerical simulation of MOS devices. Much emphasis is laid on the discussion of recent refinements to carrier transport models, e.g. drift-diffusion model, enhanced drift-diffusion equations, hydrodynamic model, and Monte Carlo simulation. Adequate models for the physical parameters are reported with suitable parameter values, e.g. carrier mobilities taking into account the various scattering mechanisms, and carrier generation-recombination including impact ionization. Examples are presented for two different types of MOS devices: on the one hand, simulation results of miniaturized MOS transistors are discussed which have been obtained by our simulator MINIMOS 5.0 with additional extensions, and on the other hand, simulation results concerning a power MOS transistor are shown which has been investigated with our device simulation program BAMBI 2.1.

Key words. simulation, MOS transistor, semiconductor equations

1. Introduction. Since 1960 when the demonstration of the practically usable MOS transistor took place [42] its development has shown to evolve dramatically. Today, about thirty years later, integrated circuits with millions of devices per single chip are manufactured. In order to minimize the number of cycles of trial and error in device fabrication improved understanding of basic device operation has attained crucial importance. Thus, numerical modeling of MOS transistors has become a basic requirement for the development of prototype devices although the first modeling attempts were made relatively late.

A two dimensional solution of the Poisson equation with application to a MOS structure probably was first published by Loeb et al. [50] and Schroeder and Muller [68] in 1968. Since then a lot of work has been contributed dealing with simulation of MOS devices due to their intrinsically two dimensional nature, e.g.: [90], [52], [38], [58], [87], [73] during the seventies, [14], [77], [88], [46], [69], [92], [55], [22] during the eighties.

Two dimensional transient simulations of MOSFET's have been carried out in, e.g.: [53], [57], [94], [89], [78], [32], [24], [43]. Three dimensional static modeling has been published in, e.g.: [7], [11], [39], [80], [49], [13], [86], [76].

The main efforts during the recent years were not focused only on transient simulations and three dimensional modeling but also on the improvement of the models of the physical parameters, e.g.: [37], [56], [91], [81], on the refinement of the mathematical model, e.g.: [67], [31], [45], on the analysis of the operation of MOSFET's at low temperature, e.g.: [82], [93], [47], [60], [79], [74], and on investigations of special physical effects, e.g.: [30], [62], [61], [28], [75], [34].

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Although most activities in MOS transistor simulation are concentrated on the analysis of ultra small devices, there also exist many publications concerning modeling of power MOSFET's, e.g.: [85], [22], [19].

2. Transport Model. The state-of-the-art in semiconductor technology enables the fabrication of semiconductor devices with characteristic lengths smaller than $1\ \mu\text{m}$ in an industrial environment. Such devices are characterized by very large electric fields and rapid spatial variations of the electric field and carrier concentrations. These variations occur over a distance comparable with the characteristic lengths of carrier transport, i.e. the average momentum relaxation length and the energy relaxation length. Owing to the very large electric fields the assumption of small perturbations around equilibrium breaks down which the drift-diffusion transport model is based on. It was realized long ago that a straightforward extension of the classical semiconductor equations would have to include the energy balance of field-driven carriers [15]. One way for this purpose is an enhanced drift-diffusion model which accounts for the energy balance. The obvious benefit of this approach is that it can be implemented without great efforts in an existing simulation program based on the basic semiconductor equations. Another practicable way to include hot carrier effects is the usage of the hydrodynamic model [65] which implies the energy conservation equation. Last not least, there is to discuss the Monte Carlo method whose most important drawback is — as usual — the enormous amount of required computing resources.

2.1. Drift-Diffusion Equations. Device Modeling based on the self-consistent solution of fundamental semiconductor equations dates back to the famous work of Gummel in 1964 [25]. Up to now, most of the worldwide used simulation tools for MOS transistors are solving the well known fundamental semiconductor equations consisting of Poisson's equation (1), continuity equations for electrons and holes (2), (3), and the current relations for both carrier types (4), (5).

$$(1) \quad \text{div}(\epsilon \cdot \text{grad } \psi) = -\rho$$

$$(2) \quad \text{div } \vec{J}_n - q \cdot \frac{\partial n}{\partial t} = q \cdot R$$

$$(3) \quad \text{div } \vec{J}_p + q \cdot \frac{\partial p}{\partial t} = -q \cdot R$$

$$(4) \quad \vec{J}_n = q \cdot \mu_n \cdot n \cdot \left(\vec{E} + \frac{1}{n} \cdot \text{grad} \left(n \cdot \frac{k_B \cdot T_n}{q} \right) \right)$$

$$(5) \quad \vec{J}_p = q \cdot \mu_p \cdot p \cdot \left(\vec{E} - \frac{1}{p} \cdot \text{grad} \left(p \cdot \frac{k_B \cdot T_p}{q} \right) \right)$$

It should be noted that the current relations in this notation already represent an extension of the classical drift-diffusion equations as there exist different temperatures T_n and T_p within the diffusion terms. The reason for this will be explained in the next section.

As it can be seen from the different activities mentioned above there are ongoing arguments in the scientific community whether these equations are adequate to describe transport in submicron devices. Particularly the current relations (4) and (5) which are the most complex equations out of the set of the basic semiconductor device equations undergo strong criticism in view of, for instance, ballistic transport [64], [35]. Their derivation from more fundamental physical principles is indeed not at all straightforward. They appear therefore with all sorts of slight variations in the specialized literature and a vast number of papers has been published where some of their subtleties are dealt with. Anyway, investigations on ultra short MOSFET's [66] do not give evidence that it is necessary to waive these well established basic equations for silicon devices down to feature sizes in the order of 0.1 microns [79].

2.2. Enhanced Drift-Diffusion Equations. In contrast of the classical drift-diffusion model the solution of the enhanced drift-diffusion equations accounts for hot-carrier transport in semiconductors. An adequate model for already existing simulation programs has been implemented first in MINIMOS [31], our well established tool for the two and three dimensional analysis of miniaturized MOS transistors.

This model is based on the work of Hänsch and Miura-Mattausch [29] which follows very closely the work of Blotekjaer [6] with the distinction that in the first place a different Ansatz for the distribution function was chosen, and secondly, and more importantly, the relaxation time approximation was not used. This provides a closed set of equations for the electron density n , the electron current density $\vec{J}_n$, energy density and energy current density. From this set of equations a self-consistent mobility and energy relaxation time have been derived rigorously [27]. This formulation covers particle as well as energy balance in a closed form. Unfortunately it would require the development of a new code to be utilized as a simulation tool. Therefore, Hänsch and Miura-Mattausch presented an approximate solution of these equations that could feasibly be used in conventional drift-diffusion codes. Their rational was based on the observation that, strictly speaking, the classical current equation is rigorously valid in the limit of small fields. High-field effects enter through the saturation of the drift velocity and are accounted for using a heuristic field-dependent mobility model, which is a local field-dependent function. Following these classical ideas, Hänsch and Miura-Mattausch introduced a local field dependent model of the mobilities μ_n and μ_p and the thermal voltages U_{T_n} and U_{T_p} to be used in the modified current relationships (4) and (5).

To model μ and U_T , however, a local solution of the general equations was performed. To this end all spatial derivatives were dropped so that the complicated system of differential equations turned into a simple algebraic equation for the energy density and in turn the mobility. This rather drastic step is justified by the observation that, in conventional simulation, mobility models are inferred from drift velocity measurements in constant electric field. The inhomogeneous situation is then accounted for by an appropriate driving force [71], [72]. Assuming infinite bulk material in the first place and replacing the electric field

by the driving force F , the electron mobility reads

$$(6) \quad \mu_n^{LISF} = \frac{2\mu_n^{LIS}}{1 + \sqrt{1 + \left(\frac{2 \cdot \mu_n^{LIS} \cdot F_n}{v_n^{sat}} \right)^2}}$$

and for the thermal voltage

$$(7) \quad U_{T_n} = \frac{k_B \cdot T_n}{q} = U_{T_0} + \frac{2}{3} \cdot \tau_n^\epsilon \cdot (v_n^{sat})^2 \cdot \left(\frac{1}{\mu_n^{LISF}} - \frac{1}{\mu_n^{LIS}} \right)$$

is obtained. Here, μ_n^{LIS} is the low-field mobility containing lattice, impurity, and surface scattering, τ_n^ϵ is the energy relaxation time, and v_n^{sat} is the saturation velocity.

Concerning the driving force F , there is some controversy about which one is the most appropriate force in the literature. The problem is that originally the local field-dependent expressions are given with $F = E$. Generalization to the inhomogeneous two-dimensional case is by no means straightforward. There are, however, two limiting cases the driving force has to obey: On the one hand it should give $F = E$ for the homogeneous situation, and on the other hand it has to give a velocity saturation for high-density gradients as well. In this situation carriers do not gain energy from the field, provided it is small, and therefore their mean velocity cannot exceed the thermal velocity, which is not necessarily the same as the high-field saturation velocity v^{sat} . Any physically motivated driving force has to fulfill these requirements. Under the present circumstances a saturation with respect to the gradient of the thermal voltage has to be included too. One appropriate driving force proposed firstly in [31] is

$$(8) \quad F_n = \left| \text{grad } \psi - \frac{1}{n} \cdot \text{grad}(n \cdot U_{T_n}) \right|.$$

2.3. Hydrodynamic Equations. The derivation of the hydrodynamic model describing the carrier transport in semiconductor devices is the Boltzmann transport equation

$$(9) \quad \frac{\partial f}{\partial t} + \vec{u} \cdot \text{grad}_x f \pm \frac{q}{m} \cdot \vec{E} \cdot \text{grad}_u f = C$$

where $f = f(x_j, u_j, t)$ is the distribution function, x_j, u_j ($j = 1, 2, 3$) are the components of the vector position $\vec{x}$ and group velocity $\vec{u}$, q is the electron charge, m denotes the effective mass of the carriers under investigation, $\vec{E}$ is the electric field, and C denotes the collision integral. In semiconductors it is necessary to write such an equation for each branch of the energy-band function, using the negative sign for electrons and the positive sign for holes. In the above written formulation the constant mass approximation is assumed, and the scalar mass m is a proper average of terms of the mass tensor relative to each extremum of the considered energy-band function. Equation (9) is a partial differential equation in the seven dimensional phase space $(\vec{x}, \vec{u}, t)$. In order to reduce this equation to a set of equations

only in the space and time coordinates usually the so called method of moments is used as described in [65].

The derivation of the hydrodynamic model implies some loss of information on the behavior of the distribution function f . The main reason for this is the absolutely necessary consideration of only a finite number of moments of the Boltzmann transport equation. Additional details can be obtained by introducing more moments, at the expense of an increasing complication of the equations. In order to obtain an equation system comparable to the drift-diffusion model, the following simplifying assumptions have to be taken: the temperature tensor is reduced to a scalar, and the so called relaxation-time approximation is used for the description of the collision terms of the intervalley and intravalley transition type.

After introducing the set of variables commonly used in semiconductor device theory one obtains a current relation

$$(10) \quad \vec{J}_n - \frac{\tau_{pn}}{q} (\vec{J}_n \cdot \text{grad}) \frac{\vec{J}_n}{n} = q \cdot \mu_n \left[\frac{k_B \cdot T_n}{q} \text{grad}(n) + n \cdot \text{grad} \left(\frac{k_B \cdot T_n}{q} - \psi \right) \right]$$

and an energy conservation equation

$$(11) \quad -\text{div} \left(\kappa_n \cdot \text{grad}(T_n) + \vec{J}_n \cdot \frac{k_B \cdot T_n}{q} \right) = \vec{J}_n \cdot \text{grad} \left(\frac{w}{q} - \psi \right) - n \cdot \frac{w - w_0}{\tau_{wn}}$$

where τ_{pn} and τ_{wn} are the momentum and energy relaxation times for electrons in the conduction band, respectively, and w_0 is the equilibrium mean energy. These two equations in connection with the continuity equation (2) whose formulation has not changed build the basic equations of the hydrodynamic transport model (for electrons). They have to be solved self-consistently with the corresponding equations for holes and with the Poisson equation.

Regarding the increased expense, the hydrodynamic model has an evident drawback compared to the drift-diffusion model: the hydrodynamic model has two further partial differential equations (the energy conservation equation for both carrier types) which have to be solved self-consistently with the other three (i.e. Poisson equation and both continuity equations). Furthermore, the formulation of the current relations is more complex than in the drift-diffusion model.

But there is no doubt about the fact that in spite of the simplifying assumptions the hydrodynamic model is more rigorous than the classical drift-diffusion model, especially concerning hot carrier transport. Let us compare the two different formulations of the current relations (4) and (10). The current relation of the classical drift-diffusion model, i.e. equation (4) with $T_n = T_p = T_L$, where T_L is the lattice temperature, can be obtained from (10) by neglecting $(\tau_{pn}/q)(\vec{J}_n \cdot \text{grad})(\vec{J}_n/n)$, which is called convective term, and the gradient of $k_B T_n/q$, which is a thermoelectric field. Furthermore, in the classical drift-diffusion model the diffusion coefficient is defined by the Einstein relation

$$(12) \quad D_{n,p} = \frac{k_B \cdot T_L}{q} \cdot \mu_{n,p}$$

whereas the analogous term in the hydrodynamic model reads

$$(13) \quad D_{n,p} = \frac{k_B \cdot T_{n,p}}{q} \cdot \mu_{n,p}.$$

If we compare now the current relations of the hydrodynamic model (10) and the enhanced drift-diffusion model (4) we see that the thermoelectric field as well as the so called generalized Einstein relation (13) are included implicitly — the only difference is the neglect of the conductive term. As reported in [65] the influence of the convective term on the solution of the partial differential equation system is very small and is, therefore, commonly neglected even in the actual simulation codes using the hydrodynamic model. Thus, from our point of view, the increased expense of the hydrodynamic equations compared to the enhanced drift-diffusion model briefly described in the previous section is not generally justified.

2.4. Hybrid Transport Model. Since the very first beginning of Monte Carlo simulations in the field of carrier transport in semiconductor devices [54] the most important drawback up to now is the enormous amount of computing resources required by this method for the solution of the Boltzmann transport equation (9). The main benefits of this method are the facts that it is simple to implement, that sophisticated physical models can be used, and that any desired physical information can easily be extracted [40].

A very promising way is the coupling of the Monte Carlo method and the drift-diffusion equations [45] which is based on the following considerations: The Monte Carlo method allows more accurate physical models and is, therefore, well suited to describe the non-equilibrium transport occurring under conditions appearing in ultra small MOS transistors (i.e. very high electric field in the active region with rapid changes over distances comparable to the mean free path of the carriers). On the other hand, for description of low field transport the drift-diffusion model which uses local transport coefficients provides sufficient accuracy. Moreover, the drift-diffusion model has turned out to be even superior to the Monte Carlo method in regions with retarded fields. Therefore, several attempts were published to combine the drift-diffusion and Monte Carlo technique in order to benefit from the different capabilities of both methods [63], [83]. In the following, a recently performed implementation of the hybrid transport model in the existing simulation tool MINIMOS shall be discussed. It should be noted that up to now this extension has been implemented only in the two-dimensional part of the program package.

Three-dimensional scattering rates in the entire device are assumed, thus neglecting quantization effects eventually occurring in an inversion layer. Multiplication of the Boltzmann transport equation with wave vector component k_i and integration of $\vec{k}$ leads to a momentum balance equation, which reads for electrons

$$(14) \quad - \left(q \cdot E_i + \frac{1}{n} \sum_{j=1}^2 \frac{\partial (n \cdot \langle \hbar \cdot k_i \cdot v_j \rangle)}{\partial x_j} \right) = \langle \int (\hbar \cdot k_i - \hbar \cdot k'_i) \cdot S(\vec{k}, \vec{k}') \cdot d^3 \vec{k}' \rangle, \quad i = 1, 2$$

where v denotes the group velocity. (It should be mentioned that in the following the index n for indicating the electrons is omitted as this carrier type is discussed in this section

only.) The average operator $\langle A \rangle$ is the mean value of $A(\vec{k})$ weighted by the local distribution function $f(\vec{x}, \vec{k})$. The left hand side can be interpreted as driving force that acts on the electron ensemble, consisting of the electric field plus diffusion term, whereas the right hand side describes the rate of momentum loss due to scatterings. This equation can be expressed in a form similar to the drift-diffusion current relation. The parameters needed in this current relation are derived in the following way.

For band structures with spheric and ellipsoidal energy surfaces the vector valued momentum loss integral is colinear with the momentum $\hbar \vec{k}$

$$(15) \quad \int (\hbar \cdot \vec{k} - \hbar \cdot \vec{k}') \cdot S(\vec{k}, \vec{k}') \cdot d^3 \vec{k}' = \hbar \cdot \vec{k} \cdot \lambda_m(E).$$

Here the proportionality factor $\lambda_m(E)$ is the momentum scattering rate. With the local average velocity the local momentum loss mobility can be defined as

$$(16) \quad \mu = q \cdot \frac{\|\langle \vec{v} \rangle\|}{\|\langle \hbar \cdot \vec{k} \cdot \lambda_m(E) \rangle\|}$$

where $\|B\|$ denotes the maximum norm of B . This definition does not rely on the relaxation time approximation and, since no effective mass occurs in this formula, extension to general bands is straightforward. In the latter case μ would have tensor property. The definition of the thermal voltage tensor (which is proportional to the temperature tensor: $U_{ij} = (k_B/q)T_{ij}$) results directly from the momentum conservation equation (14)

$$(17) \quad U_{ij} = \frac{1}{q} \cdot \langle \hbar \cdot k_i \cdot v_j \rangle.$$

This definition is independent of the underlying band structure model. Inserting these definitions in (14) one obtains a general current relation

$$(18) \quad J_i = q \cdot n \cdot \mu \cdot \left(E_i + \frac{1}{n} \cdot \sum_{j=1}^2 \frac{\partial(n \cdot U_{ij})}{\partial x_j} \right).$$

The differences between this current relation and the classical one are twofold. Firstly, the diffusion term is due to the tensor property of the thermal voltage more complicated. Secondly, the parameters μ and U_{ij} can no longer be treated simply as parameters depending on electric field or other local quantities, as it is usually done in the conventional drift-diffusion model, because they carry information of the local distribution function. By means of the Monte Carlo method these parameters are evaluated. The conventional simulator using the Monte Carlo parameters $\mu(\vec{x})$ and $U_{ij}(\vec{x})$ in the current relation (18) is then capable of recovering the Monte Carlo results for $n(\vec{x})$ and $\vec{J}_n(\vec{x})$. In this way hot electron effects, such as velocity overshoot and hot carrier diffusion, are consistently included in the conventional simulator. The solution is performed globally in the whole device, but only in the high field region mobility and temperature profiles have to be extracted from the Monte Carlo

procedure. In regions with low fields and low spatial inhomogeneities local models can be used thus saving computation time.

The continuity equation in conjunction with a drift–diffusion current relation employing a scalar temperature yields an elliptic partial differential equation which has diagonal form. If an anisotropic temperature is taken into account cross derivatives appear in the elliptic operator. Conventional device simulators solve elliptic systems which are in diagonal form, therefore a scalar temperature is desirable. Neglecting the off diagonal elements in (17) a scalar temperature is used which is the arithmetic mean value of the main diagonal elements.

For the first conduction band of silicon we use a model consisting of six anisotropic valleys with a first order correction for nonparabolicity [41]. Acoustic intravalley scattering in the elastic approximation, intervalley phonon scattering, surface roughness scattering, and coulomb scattering are taken into account. Except of the latter one all mechanisms are isotropic. In the case of surface scattering in the inversion layer the wave vector is redistributed randomly in a plane parallel to the Si–SiO₂ interface. For isotropic scattering mechanisms the momentum scattering rate does not differ from the total scattering rate. The following superposition implies independence of all scattering processes

$$(19) \quad \lambda_m(E) = \lambda_{ac}^{tot} + \lambda_{opt}^{tot} + \lambda_{surf}^{tot} + \lambda_{ion}^m.$$

Coulomb interaction is the only one to be treated separately. The momentum loss integral (15) is evaluated by the Brooks–Herring formulation for the transition rate $S(\vec{k}, \vec{k}')$ and

$$(20) \quad \lambda_{ion}^m(E) = \lambda_{ion}^{tot} - \lambda'_{ion}(E)$$

is obtained. The subtraction corresponds to the difference of initial and final wave vector in (15).

Fig. 1 depicts the energy dependencies of total (λ_{ion}^{tot}) and momentum (λ_{ion}^m) scattering rate for ionized impurities, and additionally the total scattering rates for acoustic intravalley phonons (λ_{ac}) and one representative intervalley phonon mode (emission and absorption).

3. Models for Parameters. The fundamental semiconductor equations include a set of parameters which have to be appropriately modeled in order to describe the various transport phenomena qualitatively and quantitatively correctly. The basic requirement for an estimation of most of these physical parameters is independent more or less of the underlying transport model as the Poisson equation (1) and both continuity equations (2), (3) have to be solved self-consistently in any case.

3.1. Modeling Space Charge. The Poisson equation (1) requires a model for the space charge ρ which makes use of only the dependent variables ψ, n, p and material properties. The well established approach for this model is to sum up the concentrations with the adequate charge sign multiplied with the elementary charge

$$(21) \quad \rho = q \cdot (p - n + N_D^+ - N_A^-).$$

Fig. 1: Total and momentum scattering rates used for mobility calculation.

The doping concentration is assumed usually to be fully ionized at room temperature. In order to include low temperature phenomena it seems to be necessary to include partial ionization which can be described by

$$(22) \quad N_D^+ = \frac{N_D}{1 + 2 \cdot \exp\left(\frac{E_{fn} - E_D}{k_B \cdot T}\right)}, \quad N_A^- = \frac{N_A}{1 + 4 \cdot \exp\left(\frac{E_A - E_{fp}}{k_B \cdot T}\right)}$$

in the classical way. E_D and E_A are the ionization energies of the respective donor and acceptor dopant. Typical values for $E_c - E_D$ and $E_A - E_v$ for the most common dopants in silicon are: 0.054eV for arsenic, 0.045eV for phosphorus, 0.039eV for antimony, and 0.045eV for boron. A quite complete list of values can be found in [84]. Note, that only energy differences can be given (E_c and E_v are the conduction band and the valence band, respectively). The Fermi levels E_{fn} and E_{fp} have to be appropriately related to the dependent variables by making use of Fermi statistics

$$(23) \quad n = N_c \cdot \frac{2}{\sqrt{\pi}} \cdot F_{1/2}\left(\frac{E_{fn} - E_c}{k_B \cdot T}\right), \quad p = N_v \cdot \frac{2}{\sqrt{\pi}} \cdot F_{1/2}\left(\frac{E_v - E_{fp}}{k_B \cdot T}\right).$$

N_c and N_v are the density of states in the conduction band and the valence band, respectively. The classical formulae for the density of states are given by

$$(24) \quad N_c = 2 \cdot \left(\frac{2 \cdot \pi \cdot k_B \cdot T \cdot m_n^*}{h^2} \right)^{3/2}, \quad N_v = 2 \cdot \left(\frac{2 \cdot \pi \cdot k_B \cdot T \cdot m_p^*}{h^2} \right)^{3/2}.$$

$F_{1/2}$ is the Fermi function of order 1/2 which is defined by

$$(25) \quad F_{1/2}(x) = \int_0^\infty \frac{\sqrt{y}}{1 + e^{y-x}} \cdot dy.$$

The parameters m_n^* and m_p^* are the effective masses for electrons and holes, respectively. It is worthwhile to note that the ratio of the density of states depends only on the ratio of the effective masses

$$(26) \quad \frac{N_c}{N_v} = \left(\frac{m_n^*}{m_p^*} \right)^{3/2}.$$

By means of simple algebraic manipulation with the expressions for the carrier concentrations (23) one obtains

$$(27) \quad \frac{E_{fn} - E_D}{k_B \cdot T} = G_{1/2} \left(\frac{n}{N_c} \right) + \frac{E_c - E_D}{k_B \cdot T}$$

$$(28) \quad \frac{E_A - E_{fp}}{k_B \cdot T} = G_{1/2} \left(\frac{p}{N_v} \right) + \frac{E_A - E_v}{k_B \cdot T}.$$

$G_{1/2}(x)$ is the inverse Fermi function of order 1/2 defined by

$$(29) \quad G_{1/2} \left(\frac{2}{\sqrt{\pi}} \cdot F_{1/2}(x) \right) = x.$$

It is now possible by evaluating the expressions for the density of states (24) with appropriate fits to the effective masses [23] to compute numerical values for the ionized impurity concentrations (22) using only the carrier concentrations as the dependent variables in the basic equations. However, comparisons to experiment indicate that it is better to compute the density of states from relation (26) and a fit to the intrinsic carrier concentration

$$(30) \quad n_i = \sqrt{N_c \cdot N_v} \cdot \exp \left(-\frac{E_g}{2 \cdot k_B \cdot T} \right)$$

where E_g is the band gap $E_c - E_v$ which can be modeled temperature dependent with the fit provided by Gaensslen et al. [23]. The prefactor in (30) can be fitted to experiments by

$$(31) \quad \sqrt{N_c \cdot N_v} = \exp \left\{ 45.13 + 0.75 \cdot \ln \left[\frac{m_n^*}{m_0} \cdot \frac{m_p^*}{m_0} \cdot \left(\frac{T}{300\text{K}} \right)^2 \right] \right\} \text{ cm}^{-3}.$$

With (26) and (31) it is now straightforward to compute the numerical values for the density of states. At room temperature we obtain $N_c = 5.1 \cdot 10^{19} \text{ cm}^{-3}$ and $N_v = 2.9 \cdot 10^{19} \text{ cm}^{-3}$.

It should be noted that equations (22) are valid only for moderate impurity concentrations. For heavy doping the assumption of a localized ionization energy does definitely not hold. Instead an impurity band is formed which may merge with the respective band edge, e.g. [36]. It appears to be appropriate to assume total ionization for concentrations above some threshold value and to account for a suitable functional transition between the classical formulae (22) and total ionization.

3.2. Modeling Carrier Mobilities. The next set of physical parameters to be considered carefully consists of the carrier mobilities μ_n and μ_p in the current relations (4) and (5) or (10) regarding the hydrodynamic model. The models for the carrier mobilities have to take into account a great variety of scattering mechanisms the most basic one of which is lattice scattering. The lattice mobility in pure silicon can be fitted with simple power laws

$$(32) \quad \mu_n^L = 1430 \frac{\text{cm}^2}{\text{Vs}} \cdot \left(\frac{T}{300\text{K}} \right)^{-2}, \quad \mu_p^L = 460 \frac{\text{cm}^2}{\text{Vs}} \cdot \left(\frac{T}{300\text{K}} \right)^{-2.18}.$$

The expressions (32) fit well experimental data of [1], [8], and [48].

The next effect to be considered is ionized impurity scattering. The best established procedure for this task is to take the functional form (33) of the fit provided by Caughey and Thomas [10] and use temperature dependent coefficients.

$$(33) \quad \mu_{n,p}^{LI} = \mu_{n,p}^{min} + \frac{\mu_{n,p}^L - \mu_{n,p}^{min}}{1 + \left(\frac{CI}{C_{n,p}^{ref}} \right)^{\alpha_{n,p}}}$$

$$(34) \quad \mu_n^{min} = 80 \frac{\text{cm}^2}{\text{Vs}} \cdot \left(\frac{T}{300\text{K}} \right)^{-0.45}, \quad \mu_p^{min} = 45 \frac{\text{cm}^2}{\text{Vs}} \cdot \left(\frac{T}{300\text{K}} \right)^{-0.45}$$

$$(35) \quad C_n^{ref} = 1.12 \cdot 10^{17} \text{cm}^{-3} \cdot \left(\frac{T}{300\text{K}} \right)^{3.2}$$

$$C_p^{ref} = 2.23 \cdot 10^{17} \text{cm}^{-3} \cdot \left(\frac{T}{300\text{K}} \right)^{3.2}$$

$$(36) \quad \alpha_{n,p} = 0.72 \cdot \left(\frac{T}{300\text{K}} \right)^{0.065}$$

The fits (34)-(36) are from [33]. Similar data have been provided in [3] and [20].

In view of partial ionization one should consider neutral impurity scattering. However, in view of the uncertainty of the quantitative values for ionized impurity scattering it seems not to be worthwhile to introduce another scattering mechanism with additional fitting parameters.

Particular emphasis has to be put on surface scattering which is modeled with an expression suggested by Seavey [70]

$$(37) \quad \mu_{n,p}^{LIS} = \frac{\mu_{n,p}^{ref} + (\mu_{n,p}^{LI} - \mu_{n,p}^{ref}) \cdot (1 - F(y))}{1 + F(y) \cdot \left(\frac{S_{n,p}}{S_{n,p}^{ref}} \right)^{\gamma_{n,p}}}$$

with

$$\begin{aligned} \mu_n^{ref} &= 638 \frac{\text{cm}^2}{\text{Vs}} \cdot \left(\frac{T}{300\text{K}} \right)^{-1.19}, & \mu_p^{ref} &= 160 \frac{\text{cm}^2}{\text{Vs}} \cdot \left(\frac{T}{300\text{K}} \right)^{-1.09} \\ F(y) &= \frac{2 \cdot \exp \left(- \left(\frac{y}{y^{ref}} \right)^2 \right)}{1 + \exp \left(-2 \cdot \left(\frac{y}{y^{ref}} \right)^2 \right)} \\ S_n &= \max \left(0, \frac{\partial \psi}{\partial y} \right), & S_p &= \max \left(0, -\frac{\partial \psi}{\partial y} \right). \end{aligned}$$

S_n^{ref} is assumed to be $7 \cdot 10^5 \frac{\text{V}}{\text{cm}}$, S_p^{ref} is $2.7 \cdot 10^5 \frac{\text{V}}{\text{cm}}$, γ_n is 1.69, γ_p is 1.0, and y^{ref} is 10 nm. The formulae for surface scattering are definitely not the ultimate expressions. They just fit quite reasonably experimental observations.

Velocity saturation can be modeled with formulae (38). The expression for the electron mobility comes directly from the derivation of the enhanced drift-diffusion model (6). For the hole mobility the same functional form has been assumed and has been fitted to experimental data.

$$(38) \quad \mu_n^{LISF} = \frac{2\mu_n^{LIS}}{1 + \sqrt{1 + \left(\frac{2 \cdot \mu_n^{LIS} \cdot F_n}{v_n^{sat}} \right)^2}}, \quad \mu_p^{LISF} = \frac{\mu_p^{LIS}}{1 + \frac{\mu_p^{LIS} \cdot F_p}{v_p^{sat}}}$$

F_n and F_p are the effective driving forces which read according to (8)

$$(39) \quad F_n = \left| \text{grad } \psi - \frac{1}{n} \cdot \text{grad}(n \cdot U_{Tn}) \right|, \quad F_p = \left| \text{grad } \psi + \frac{1}{p} \cdot \text{grad}(p \cdot U_{Tp}) \right|.$$

The used saturation velocities are given by

$$(40) \quad \begin{aligned} v_n^{sat} &= 1.45 \cdot 10^7 \frac{\text{cm}}{\text{s}} \cdot \sqrt{\tanh \left(\frac{155\text{K}}{T} \right)} \\ v_p^{sat} &= 9.05 \cdot 10^6 \frac{\text{cm}}{\text{s}} \cdot \sqrt{\tanh \left(\frac{312\text{K}}{T} \right)}. \end{aligned}$$

The functional form of these fits is after [1], and the experimental data matched are [1], [9], and [8]. An eventual dependence on the crystallographic orientation which one would deduce from [2], [44] is presently not taken into account.

3.3. Modeling Carrier Temperatures. To describe carrier heating properly one has to account for local carrier temperatures $T_{n,p}$ in the current relations. As already discussed in the previous chapters this can be achieved by either solving the hydrodynamic equations, or by using a model obtained by series expansions of the solution to the energy conservation equations, e.g. by using the enhanced drift-diffusion model. As already discussed we believe that the latter is sufficient for silicon devices. Applying equation (7) also for holes the enhanced drift-diffusion model yields for the electronic voltages

$$(41) \quad U_{T_{n,p}} = \frac{k_B \cdot T_{n,p}}{q} = U_{T_0} + \frac{2}{3} \cdot \tau_{n,p}^\epsilon \cdot (v_{n,p}^{sat})^2 \cdot \left(\frac{1}{\mu_{n,p}^{LISF}} - \frac{1}{\mu_{n,p}^{LIS}} \right).$$

The energy relaxation times $\tau_{n,p}^\epsilon$ are in the order of 0.5 picoseconds. They should be modeled as functions of the local doping concentration as motivated by the following reasoning: The product of carrier mobility times electronic voltage which symbolizes a diffusion coefficient must be a decreasing function with increasing carrier voltage [5]. Its maximum is attained at thermal equilibrium. Therefore, the relation

$$(42) \quad \mu_{n,p}^{LISF} \cdot U_{T_{n,p}} \leq \mu_{n,p}^{LIS} \cdot U_{T_0}$$

must hold. Substituting (41) into (42) and rearranging terms one obtains for the energy relaxation times

$$(43) \quad \tau_{n,p}^\epsilon \leq \frac{3}{2} \cdot U_{T_0} \cdot \frac{\mu_{n,p}^{LIS}}{(v_{n,p}^{sat})^2}.$$

In MINIMOS the energy relaxation times are modeled on this basis with a fudge factor γ in the range $[0, 1]$ and a default value of 0.8

$$(44) \quad \tau_{n,p}^\epsilon = \gamma \cdot \frac{3}{2} \cdot U_{T_0} \cdot \frac{\mu_{n,p}^{LIS}}{(v_{n,p}^{sat})^2}.$$

For vanishing doping one obtains the maximum energy relaxation times which are at room temperature $\tau_n^\epsilon = 4.44 \cdot 10^{-13}$ s, $\tau_p^\epsilon = 2.24 \cdot 10^{-13}$ s.

3.4. Modeling Carrier Generation/Recombination. The right hand sides of the continuity equations are represented by the carrier generation/recombination term which consists of various mechanisms.

An adequate model for thermal generation/recombination is given by the well known Shockley-Read-Hall term [71]

$$(45) \quad R^{SRH} = \frac{n \cdot p - n_i^2}{\tau_p \cdot (n + n_1) + \tau_n \cdot (p + p_1)}$$

with

$$(46) \quad n_1 = n_0 \cdot \frac{1 - f_{t0}}{f_{t0}}, \quad p_1 = p_0 \cdot \frac{f_{t0}}{1 - f_{t0}}$$

and the doping dependent carrier life times

$$(47) \quad \tau_n = \frac{\tau_{n0}}{1 + \frac{N_D + N_A}{N_n^{ref}}}, \quad \tau_p = \frac{\tau_{p0}}{1 + \frac{N_D + N_A}{N_p^{ref}}}.$$

f_{t0} in equation (46) denotes the fraction of occupied traps at equilibrium, n_0 and p_0 are the equilibrium carrier densities. The doping dependent empirical model for the carrier life times (47) accounts for additional generation/recombination centers at high doping. The parameters are chosen to fit experimental findings, e.g. $\tau_{n0} = 3.94 \cdot 10^{-4}$ s, $N_n^{ref} = 7.1 \cdot 10^{15} \text{ cm}^{-3}$, $\tau_{p0} = 3.94 \cdot 10^{-5}$ s, $N_p^{ref} = 7.1 \cdot 10^{15} \text{ cm}^{-3}$ [18].

Auger recombination is modeled by

$$(48) \quad R^{AU} = (C_{cn} \cdot n + C_{cp} \cdot p) \cdot (n \cdot p - n_i^2)$$

with appropriate temperature dependent Auger coefficients [21]

$$(49) \quad \begin{aligned} C_{cn} &= 2.8 \cdot 10^{-31} \frac{\text{cm}^6}{\text{s}} \cdot \left(\frac{T}{300\text{K}} \right)^{0.14} \\ C_{cp} &= 9.9 \cdot 10^{-32} \frac{\text{cm}^6}{\text{s}} \cdot \left(\frac{T}{300\text{K}} \right)^{0.2} \end{aligned}$$

As model for impact ionization the old Chynoweth formulation (50) can be used. Although the actuality of this model [12] has been discussed heavily in the scientific community, it still seems that it fulfills the requirements for device simulation quite satisfactorily — at least not less than other models published in literature.

$$(50) \quad R^{II} = -\alpha_n \cdot \frac{|\vec{J}_n|}{q} - \alpha_p \cdot \frac{|\vec{J}_p|}{q}$$

$$(51) \quad \alpha_{n,p} = \alpha_{n,p}^{\infty} \cdot \exp\left(-\frac{\beta_{n,p}}{E}\right)$$

The coefficients of (51) can be modeled temperature dependent by (52) and (53) to fit experimental data [16], [17], [59].

$$(52) \quad \begin{aligned} \alpha_n^{\infty} &= 7 \cdot 10^5 \text{ cm}^{-1} \cdot \left(0.57 + 0.43 \cdot \left(\frac{T}{300\text{K}} \right)^2 \right) \\ \alpha_p^{\infty} &= 1.58 \cdot 10^6 \text{ cm}^{-1} \cdot \left(0.58 + 0.42 \cdot \left(\frac{T}{300\text{K}} \right)^2 \right) \end{aligned}$$

$$\begin{aligned}
(53) \quad \beta_n &= 1.23 \cdot 10^6 \frac{\text{V}}{\text{cm}} \cdot \left(0.625 + 0.375 \cdot \left(\frac{T}{300\text{K}} \right) \right) \\
\beta_p &= 2.04 \cdot 10^6 \frac{\text{V}}{\text{cm}} \cdot \left(0.67 + 0.33 \cdot \left(\frac{T}{300\text{K}} \right) \right)
\end{aligned}$$

A comparison to other models for impact ionization can be found in [4]. The scatter among the results of the different models indicates that more effort is necessary in this area to better understand impact ionization and to obtain a more rigorous description of this phenomenon.

4. Simulation results. This section deals with recently performed work concerning MOS transistor simulation within the Institute for Microelectronics at the Technical University of Vienna. Two of these topics are new and not yet released implementations in the program package MINIMOS [86], and a third work is occupied with the simulation of power MOSFET's with the more universal usable simulator BAMBI [22].

4.1. The Hybrid Model applied to a MOSFET. The coupling of Monte Carlo method and drift-diffusion model briefly described in section 2.4 has been implemented in MINIMOS in order to justify the efficiency and accuracy of that compound model with respect to ultra small MOS transistors [45]. An n-channel MOSFET with $L_{\text{gate}} = 0.25 \mu\text{m}$, $t_{\text{ox}} = 5 \text{ nm}$ was simulated at room temperature using the combined technique. The device has a metallurgical channel length of $L_{\text{eff}} = 0.15 \mu\text{m}$ and exhibits a threshold voltage $U_t = 0.23 \text{ V}$.

For practical simulation of a MOSFET, electrons are injected in source, where they fully thermalize before entering the channel [67]. In the region of interest, usually near the drain, a sufficient large number of particles is supplied by a particle split algorithm, thus reducing the statistical uncertainty of the results. The bias conditions for the following results are $U_{\text{GS}} = U_{\text{DS}} = 2.5 \text{ V}$ and $U_{\text{BS}} = 0 \text{ V}$.

The two-dimensional distribution of the main diagonal temperatures in the area near the drain are depicted in Fig. 2. The lateral temperature T_{xx} has a maximum value at the surface, while the maximum of T_{yy} is shifted away from the surface. Degradation due to hot electron injection into the oxide can be more accurately modeled by using the spatial distribution of T_{yy} than by using the scalar temperature obtained from average energy. In the performed simulations the off-diagonal temperatures never exceeded 15% of the main-diagonal elements.

Comparing the left and the right part of Fig. 3 we see that the large normal field within the inversion layer does not appear in the driving force. This is obvious since the normal field is compensated by diffusion. The electric field in the source junction is compensated by diffusion as well, since in this area the driving force calculated by Monte Carlo vanishes. The field peak near the drain edge however appears almost unchanged in the driving force, thus accelerating and heating up the electron gas in this area.

Fig. 2: Lateral and transversal temperatures in a quarter micron MOSFET (units [1000 K]).

Velocity of electrons in the channel is plotted in Fig. 4. The effective channel extends from $0.05\text{ }\mu\text{m}$ to $0.2\text{ }\mu\text{m}$, the positions of the junctions of source and drain subdiffusion, respectively. In the first half of this range the surface velocity (curve A) is lower than the velocities within the inversion layer since the electrons are pressed towards the surface. Near drain the pressing force has opposite direction and the electron velocity is maximum at the surface. The field peak near drain induces a velocity overshoot of 90% related to the bulk saturation velocity. A comparison of the two parts of Fig. 3 has shown that in the overshoot region diffusion is not important. Therefore, velocity overshoot is treated in this model more like a drift phenomenon. It is reproduced by the drift-diffusion current relation (18) by incorporating the nonlocal mobility which will also be increased compared to a local mobility, which cannot produce velocities larger than the bulk saturation velocity.

4.2. Non-Degenerate Gate Effect in MOSFET's. Implanted gate MOS devices have become common in submicron technologies. For that reason the simulator MINIMOS has been extended to solve self-consistently the basic semiconductor equations also in the polygate area (fully including non-planar devices) [26]. The Poisson equation is solved in the total simulation area (from y_t until y_B in Fig. 5). For the continuity equations two approaches have been implemented. In the first one, both discretized continuity equations are solved in the poly-gate simultaneously with the bulk area (from y_G until y_B in Fig. 5). This approach is interesting for the transient simulation, however a proper modeling of mobility and generation-recombination phenomena in polysilicon (e.g. grain boundary recombination) is necessary. The investigations presented here are restricted to steady-state conditions. Then the poly-gate is in thermodynamic equilibrium (net recombination vanishes and leakage currents are negligible). A unique and constant Fermi level exists in the poly-gate, which

Fig. 3: Electric field (left) and driving force (right) calculated by Monte Carlo in the same device.

enables the carrier concentrations to be calculated analytically as a function of the local potential ψ . In such an approach the discretization error of the continuity equation (specially at the non-planar gate/oxide interface — see Fig. 5) is avoided. This permits that band gap narrowing and Fermi-Dirac statistics can be implemented in a simpler way than in the first. For a detailed description of the model of heavy doping effects in the gate the interested reader is referred to [26].

Investigations have been shown that there is a strong influence of the charge at the polysilicon/oxide interface on the field penetration into the gate, and therefore on the flat-band and threshold voltage. There is not much information about the nature of this charge in the literature. Since the polysilicon is deposited over oxide, it is believed that the gate/oxide interface is worse than the interface between thermally grown oxide and bulk silicon. In the actually performed simulations there is a fixed oxide and interface trapped charge (with both donor and acceptor nature) incorporated. The traps at the grain boundaries in the polysilicon presently have not been taken into account. If the doping is several times higher than the equivalent volume trap density in polysilicon N_{tvol} (surface trap density at grain boundary/grain size), the trapped charge is negligible compared to the space charge due to impurity ions.

In order to investigate the impact of the poly-gate depletion on the characteristics of thin oxide submicron MOSFET's quarter micron planar devices have been analyzed. The devices have 5 nm oxide thickness, threshold voltage ± 0.25 V, and are designed for room temperature operation. Multiple implanted source-drain profiles (Fig. 5) are reconstructed from the data in literature. For the p-gate device it is assumed that there is no boron penetration. Fig. 6 shows the distribution of the potential, field and electron and hole

Fig. 4: Electron average velocity along the channel. Curve A: at the Si-SiO₂ interface. Curves B and C: 5 nm and 10 nm away from interface.

concentrations in the gate of the n-channel/n-gate device with the specifications $t_{ox} = 5$ nm, $L = 0.25$ μ m, $N_g = 10^{19}$ cm⁻³, and the bias conditions $U_{DS} = U_{GS} = 2$ V. Due to thin oxide, medium ionized impurity concentration at the gate/oxide interface (10^{19} cm⁻³ and high gate bias (2 V), a remarkable potential drop occurs in the gate as can be seen in the upper left (electric potential) and lower left drawing (electric field) of Fig. 6. Note that the view of the upper pictures is another than the view of the lower drawings in Fig. 6. These different representations were necessary in order to make the distributions more illustrative. The distributions of the carrier densities (upper right: electrons, lower right: holes) are plotted in logarithmic scale.

The gate-drive is reduced about 20% at the channel end near the source. Note that for

Fig. 5: Simulation area of the polysilicon-gate device.

this device the inversion in the polygate takes place (beginning at the source channel–end) at $U_{GS} \approx 3.7\text{ V}$, leading to the recovery of the transconductance (experimental finding in [51]).

The threshold voltage and the potential drop in poly–gate at the threshold versus ionized impurity concentration near the gate/oxide interface N_g are shown for the p–gate/p–channel device in Fig. 7. The charge at the gate/oxide interface Q_{go} which is the parameter of the curves in Fig. 7 (here assumed as fixed) has a strong influence on the voltage drop in the poly–gate and therefore on the threshold voltage. Assuming Q_{go} to be a positive charge the voltage drop in the gate is increased for a p–gate/p–channel device, while in an n–gate/n–channel device a positive Q_{go} has a screening effect.

The fall–off of the drain current in the saturation region with N_g as parameter is shown in Fig. 8. For common values Q_{go} has minor influence, and N_g is the main parameter in determination of the drain current degradation. In order to suppress totally the reduction of the gate drive, the activated impurity concentration near the gate/oxide interface must be at least $4 \cdot 10^{19}\text{ cm}^{-3}$ for the analyzed 5 nm oxide devices.

The relative ratio of the effective and the terminal gate–source voltage is given roughly by (assuming solely depletion in the gate)

$$(54) \quad \frac{2}{1 + \sqrt{1 + \frac{2 \cdot \epsilon_{ox}^2 \cdot U_{GS}}{q \cdot \epsilon_{pg} \cdot t_{ox}^2 \cdot N_g}}}$$

where ϵ_{ox} and ϵ_{pg} are the permittivities in oxide and polysilicon–gate, respectively. Applying different scaling rules on t_{ox} and U_{GS} the poly–gate effect becomes more or less severe by miniaturization. For instance, for the device of Fig. 8 the reduction of the current at the gate and drain supply voltage of -2 V is 20% at $N_g = 10^{19}\text{ cm}^{-3}$, while for a simulated 10 nm–oxide ($0.5\text{ }\mu\text{m}$) device the corresponding reduction was 12% at -5 V .

4.3. Simulation of power MOS transistors. Power devices are characterized by large and often complex geometries and extreme bias conditions. Therefore, simulation tools applicable to the analysis of power devices have to fulfill partially other requirements than simulators for ultra small semiconductor devices. In this section, a typical example for a power MOS transistor simulation performed at our institute will be presented. The simulation has been carried out with the program package BAMBI, a “Basic Analyzer for MOS and Bipolar devices” [22].

Fig. 9 shows the geometry of an investigated DMOS–structure. The dimensions of that geometry are identical to the used simulation area. Source, gate, drain, and bulk contacts are indicated by S, G, D, and B respectively. Note the very small highly n–doped source regions as well as the small step in the source–gate area of approximately $1\text{ }\mu\text{m}$ (small compared to the dimensions of the whole device amounting $25\text{ }\mu\text{m} \times 55\text{ }\mu\text{m}$). The drain region starts in a depth of approximately $6\text{ }\mu\text{m}$ and is connected via a highly n–doped implant to the drain

contact at the surface. The bulk contact exists only in this “test transistor” for simplification while in a real device the substrate is contacted at the surface as well.

The net doping concentration of the investigated device is shown in Fig. 10 in a logarithmic scale. The viewing direction for both drawings is given by the lower left edge of the geometry definition within Fig. 9. In the left plot of Fig. 10 which represents the doping concentration over the whole simulation area one can see the wide drain region in the depth which is connected to the surface. The step at the left vertical axis is coming from the different doping types (donors below the surface and acceptors in the depth as to be seen in Fig. 9).

The right drawing of Fig. 10 gives a detailed view of the doping profile in the channel and source region which is indicated in the left plot by two peaks only. The two peaks in the right drawing are lying at the surface and represent the source regions which are isolated by a highly doped p-implant. If one follows the doping profile at the surface one sees (from left to right) the n-doped substrate, the np-junction to the channel, and finally the pn^+ -junction between channel and source. It should be noted that the doping of this transistor has been approximated simply by superposition of Gaussian profiles.

For the calculation of distributions of internal physical quantities the bias conditions $U_{DS} = 75\text{ V}$, $U_{GS} = 20\text{ V}$, and $U_{BS} = 0\text{ V}$ were used. Fig. 11 shows the electric potential distribution (left part) and the electron density in a logarithmic scale (right part). Note that the viewing directions are not the same: while in the left drawing the front point (with the coordinates 0 lateral/30 vertical) corresponds to the right edge at the surface in Fig. 9 (drain contact) the front point of the right drawing is located $5\text{ }\mu\text{m}$ below the center of the surface. The right axis (“lateral”) of the potential distribution represents, therefore, the surface where the source region with zero voltage can easily be located. The gate-source voltage of 20 V is not visible at first glance. It is indicated only by a slight kink in the contour of the potential distribution near the source area. Following the potential from the source in vertical direction it reaches the applied drain voltage (75 V) and decreases again to the value of the substrate contact (0 V).

As mentioned before the view of the electron density distribution (right drawing of Fig. 11) is different: the left axis (“lateral”) is parallel to the surface but in a depth of $5\text{ }\mu\text{m}$ whereas the right axis (“vertical”) matches the origin. As it can be seen from the simulation geometry Fig. 9 the origin is located outside the device, thus, in that region the electron concentration is vanishing. (It should be mentioned that the logarithmic representation is dropped at zero value.) Starting the path of the electrons at the source one can nicely observe the high concentration in the channel and the bending of the electron flow into the depth.

Fig. 6: Potential, electric field, electron and hole distributions in an n-gate/n-channel device.

Fig. 7: Threshold voltage of p-gate/p-channel device.

Fig. 8: Transfer characteristics in the saturation of p-gate/p-channel device.

Fig. 9: Simulation area of the investigated DMOS-structure.

Fig. 10: Net doping density in the whole simulation area (left) and in the channel region (right).

Fig. 11: Electric potential (left) and electron density (right) distribution.

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