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Nonlinear hole transport and nonequilibrium thermodynamics in group III-nitrides under the influence of electric fields

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A theoretical study on the nonlinear transport of holes and of the nonequilibrium thermodynamic characteristics of *p*-doped wurtzite gallium nitride (GaN), aluminium nitride (AlN), and indium nitride (InN), under the influence of moderate to high electric fields, is presented. It is based on a nonlinear quantum kinetic theory which provides a description of the dissipative phenomena developing in the system. The ultrafast time evolution of the hole drift velocity and of the quasitemperatures of holes and longitudinal optical phonons are obtained. The steady state is analyzed by determining the dependence on the electric field of the nonequilibrium thermodynamic state and of the non-Ohmic mobility. A velocity overshoot is evidenced. © 2007 American Institute of Physics. [DOI: 10.1063/1.2785976]

I. INTRODUCTION

Nitride semiconductors such as gallium nitride (GaN), aluminium nitride (AlN), and indium nitride (InN) have been the object of intense research in recent years, a consequence of the large technological interest associated with their applications in blue/UV light-emitting diodes and diode lasers (see, for example, Refs. 1–4). Promising characteristics for power-field-effect devices and high-performance and high-frequency transistors,⁵ and the lack of information concerning some of the bulk III-nitrides' basic properties have led to the realization of several recent studies on their high-field steady-state transport properties.^{6–16} However, only a few of them were concerned with the transient transport regime,^{17–19} whose understanding is important for the improvement of nitride-based devices with submicron channels and high cut-off frequency.

The bulk nitride transport properties have been calculated using Monte Carlo simulations and Boltzmann transport equations. Those results should be recognized as approximations since they are highly dependent on the details of the band structure parameters which are as yet not well known. The transient transport regime in the III-nitrides follows in a subpicosecond scale (<0.3 ps), during which an overshoot in the electron velocity can be evidenced for high enough electric fields.^{18–20} These are studies related to *n*-doped III-N that we have considered in several papers.^{21–25}

Seeking a better understanding of the same characteristics in *p*-doped wurtzite GaN, AlN, and InN, we perform a theoretical study of their transient and steady-state transport properties resorting to a powerful, concise, and soundly based kinetic theory for far from equilibrium systems.²⁶ It is the one founded on a nonequilibrium statistical ensemble

formalism (NESEF),^{27,28} which provides an elegant, practical, and physically clear picture for describing irreversible processes,²⁹ as, for example, in semiconductors far from equilibrium,²⁰ which is the case considered here. An alternative treatment of extended use nowadays consists of computational methods created to model physical systems at the microscopic level, in time and distance. This began with nonequilibrium molecular dynamics (NMD) which together with the so-called Monte Carlo methods (e.g., Ref. 30) and belong to the specialization of Computational Physics.³¹ The above-mentioned NESEF-based nonlinear quantum kinetic theory provides, as said, an excellent physical picture without resort to modeling, and it appears that both approaches provide comparable results in specific cases, as shown, for example, in Ref. 24.

In the present work, through the numerical solution of relevant nonlinear quantum transport equations based on NESEF, we have characterized the hole drift velocity, mobility, and quasitemperatures (of holes and phonons) in wurtzite *p*-doped GaN, AlN, and InN. In these calculations, the interaction between phonons and carriers are considered, namely, the polar optical scattering (Fröhlich interactions with Lo phonons), the deformation and piezoelectric potentials in the case of acoustic phonons, and also the scattering by ionized impurities.

II. THE EVOLUTION EQUATIONS

To investigate the transport transient and steady-state behavior in *p*-doped AlN, GaN, and InN, we derive in the NESEF-based nonlinear quantum kinetic theory^{26,28} the evolution equations for the energy of holes and longitudinal and acoustic phonons, $E_h(t)$, $E_{Lo}(t)$, $E_{ac}(t)$, and the hole momentum $\mathbf{P}_h(t)$ (along similar lines as already applied to the study of electron transport phenomena in *n*-doped III-nitrides^{21–25}). Associated with the above quantities are five intensive non-

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equilibrium thermodynamic variables (the Lagrange multipliers that NESEF introduces), which are interpreted as related to time-evolving nonequilibrium temperatures (referred to as quasitemperatures) for carriers and longitudinal optical and acoustical phonons, which we call $T_h^*(t)$, $T_{Lo}^*(t)$, and $T_{ac}^*(t)$; the hole drift velocity $\mathbf{v}_h(t)$; and a nonequilibrium chemical potential $\mu^*(t)$ (quasichemical potential).^{28,29,32,33} It should be noticed that such choice is based on the following: (1) The carriers are internally thermalized, being brought to such condition by the efficient action of the Coulomb interaction among them, together with the collisions with Lo phonons, a randomization in energy that follows in a pico-to subpicosecond scale;³⁴ (2) Similarly, the linear momentum is also rapidly randomized presenting a uniform shift generated by the presence of the constant electric field; this can be seen in studies with the use of the present formalism, when the results are compared with experimental data and with calculations in the nonequilibrium molecular dynamics method;²⁴ (3) In the case of the Lo phonons, we recall that they do not present a randomized energy, but a certain inhomogeneous distribution in reciprocal space. This inhomogeneity in the case of electrons as carriers occurs in a very small region of the Brillouin zone (typically of the order of 10^{-14} of its whole volume of it),³⁵ and it is less pronounced in the case of holes which present a larger mass and weaker polar interaction than electrons, and such effect can be neglected in the present study.

Resorting to the use of the nonlinear quantum kinetic theory based on the nonequilibrium ensemble formalism (in which, we recall, the kinetic equations are Heisenberg equations of motion of the basic dynamical variables averaged over the nonequilibrium ensemble) in the Markovian approximation which is satisfactory in the present case³⁶ (we recall that the resulting collision integrals take then the form of the golden rule of quantum mechanics averaged over the nonequilibrium ensemble), it follows that³⁷

$$\frac{d}{dt}E_h(t) = \frac{e\mathbf{E}}{m_h} \cdot \mathbf{P}_h(t) - J_{E_h}^{(2)}(t), \quad (1)$$

$$\frac{d}{dt}\mathbf{P}_h(t) = p e \mathbf{E} - \mathbf{J}_{\mathbf{P}_h}^{(2)}(t), \quad (2)$$

$$\frac{d}{dt}E_{Lo}(t) = J_{E_{Lo}}^{(2)}(t) - J_{Lo,an}^{(2)}(t), \quad (3)$$

$$\frac{d}{dt}E_{ac}(t) = J_{E_{ac}}^{(2)}(t) + J_{Lo,an}^{(2)}(t) - J_{ac,dif}^{(2)}(t), \quad (4)$$

where p is the hole concentration, m_h^* is the hole effective mass, and \mathbf{E} is the applied electric field. The first term on the right of Eq. (1) stands for the rate of energy that the external electric field \mathbf{E} (of intensity \mathcal{E}) transfers to the carriers, while the second is the one taking account of the transfer of energy to the phonons. The first term on the right of Eq. (2) is the drifting force provided by the electric field, and the second is the rate of change of the carrier's momentum as a result of collision with phonons and impurities. In Eq. (3), we have at the right, first, the rate of change of the Lo-phonon energy

due to their interaction with the carriers; the last term accounts for the relaxation to the acoustic phonons due to anharmonic interactions. In Eq. (4), similarly to Eq. (3), the first term $J_{E_{ac}}^{(2)}(t)$ is the rate of change of the ac phonon energy due to their interaction with the carriers; the second one is the relaxation due to anharmonic interaction with the Lo phonons, and the last one is the contribution of thermal diffusion to the reservoir (for the case of carriers and Lo phonons this latter type of contribution is very small and has been neglected).

The detailed expressions for the collision operators are given in Refs. 21 and 37. It should be noticed that in Eqs. (1)–(4) the left-hand sides contain the basic variables: $E_h(t)$, $E_{Lo}(t)$, $E_{ac}(t)$, $\mathbf{P}_h(t)$, but the collision integrals on the right-hand sides depend on the nonequilibrium thermodynamic variables: $T_h^*(t)$, $T_{Lo}^*(t)$, $T_{ac}^*(t)$, $\mu^*(t)$, $\mathbf{v}_h(t)$; therefore, to close the system of equations of evolution, the relation between both types of variables, which consists in the coupled set of the so-called nonequilibrium thermodynamic equations of state,^{25,27–29,32,33} must be established. They are

$$E_h(t)/V = p \left[\frac{3}{2} k_B T_h^*(t) + \frac{1}{2} m_h^* v_h^2(t) \right], \quad (5)$$

$$\mathbf{P}_h(t)/V = p m_h^* \mathbf{v}_h(t), \quad (6)$$

$$E_{Lo}(t)/V = \frac{\hbar \omega_{Lo}}{V_{cell}} \nu_{Lo}(t), \quad (7)$$

$$E_{ac}(t)/V = \frac{3k_B T_{ac}^*(t)}{V_{cell}}, \quad (8)$$

where V is the volume, V_{cell} is the volume of the unit cell, ω_{Lo} is the dispersionless frequency (Einstein model) of the Lo phonons, $\nu_{Lo}(t)$ is the population of the latter at temperature T_{Lo}^* , and k_B is the Boltzmann constant.

The solution of the set of equations of evolution, coupled with the set of nonequilibrium equations of state, proceeds as in the case of n -doped materials, which is presented in Refs. 21–25. For a given concentration, the dependence of the hole's quasitemperature, quasichemical potential, and drift velocity, as well as the evolution of the phonon's quasitemperature, was derived. It needs to be noticed that the formalism leads, in a closed calculation, to a hole population of the form

$$f_k(t) = (\exp\{\beta_h^*(t)[\epsilon_k - \mu_h^*(t) - \mathbf{v}_h(t) \cdot \hbar \mathbf{k}]\} + 1)^{-1} \quad (9)$$

(which has acquired the form of a time-dependent—i.e., changing in time with the evolution of the nonequilibrium macroscopic state of the system—drifted Fermi-Dirac-like distribution; this is a consequence of the internal thermalization of the electrons as a result of Coulomb interaction (as already noticed). However, in all previous calculations and also in the present one, it has been used as a statistical non-degeneratelike limit for the determination of the (nonequilibrium) thermodynamic state of the holes, meaning that, in Eq. (9), 1 can be neglected compared with the exponential term. This is a satisfactory approximation once the corresponding criterion that $p\lambda_T^3 \ll 1$ in all cases is verified, where p is the

concentration and $\lambda_T(t) = \hbar / \sqrt{m_h^* k_B T_h^*(t)}$ is the de Broglie nonequilibrium-thermal wavelength, which is the situation we are analyzing here.

In the momentum evolution equation, Eq. (2), it results that the collision integrals attain the form

$$-\mathbf{J}_{\mathbf{P}_h}^{(2)}(t) = \tau_{\mathbf{P}_h}^{-1}(t, \mathcal{E}) \mathbf{P}_h(t), \quad (10)$$

where $\tau_{\mathbf{P}_h}(t)$ is a relaxation time which takes the form of Mathiessen's rule,³⁸ namely,

$$\tau_{\mathbf{P}_h}^{-1}(t, \mathcal{E}) = \tau_{\mathbf{P}_h, \text{LO}}^{-1}(t, \mathcal{E}) + \tau_{\mathbf{P}_h, \text{ac}}^{-1}(t, \mathcal{E}) + \tau_{\mathbf{P}_h, \text{imp}}^{-1}(t, \mathcal{E}), \quad (11)$$

involving the relaxation times associated with collisions with the Lo (via Fröhlich interaction) and ac phonons (via deformation and piezoelectric potentials) and with impurities.

It can be noticed that, because of Eq. (10), we can write Eq. (2) as

$$\frac{1}{V} \frac{d}{dt} \mathbf{P}_h(t) = pe\mathbf{E} - \frac{\mathbf{P}_h(t)}{V \tau_{\mathbf{P}_h}(t, \mathcal{E})}, \quad (12)$$

which can be integrated to obtain that

$$\frac{1}{V} \mathbf{P}_h(t) = pe \tau_{\text{ch}}(\mathcal{E}; t) \mathbf{E}, \quad (13)$$

where the characteristic time τ_{ch} is given by

$$\tau_{\text{ch}}(\mathcal{E}; t) = \frac{\int_0^t dt' \exp \left\{ \int_0^{t'} dt'' \tau_{\mathbf{P}_h}^{-1}(t'', \mathcal{E}) \right\}}{\exp \left\{ \int_0^t dt' \tau_{\mathbf{P}_h}^{-1}(t', \mathcal{E}) \right\}}, \quad (14)$$

where, we stress, the momentum relaxation time $\tau_{\mathbf{P}_h}(t, \mathcal{E})$ depends on the instantaneous nonequilibrium macroscopic state of the system, and thus changes in time as the latter does and we recall that the dependence on the electric field strength \mathcal{E} is indirect, that is, through the dependence of the nonequilibrium thermodynamic variables T^* and v . However, after a very rapid transient (in the picosecond scale) a *steady state* sets in, and then for times t much larger than the transient time, the characteristic time of Eq. (13) coincides with the momentum relaxation time, see Fig. 1 where the evolution of the ratio $r(t) = \tau_{\text{ch}}(t) / \tau_{\mathbf{P}_h}(t)$ is shown.

Let us consider the steady state where the current density \mathbf{j} is then [cf. Eqs. (2) and (3)] given by

$$\mathbf{j} = ep \mathbf{P}_h / m_h^* = p(e^2 / m_h^*) \tau_{\mathbf{P}_h}(\mathcal{E}) \mathbf{E} \equiv \sigma_h(\mathcal{E}) \mathbf{E}, \quad (15)$$

which is a Drude-type expression, but with a conductivity, dependent on the electric field, given by

$$\sigma_h(\mathcal{E}) = (e^2 / m_h^*) p \tau_{\mathbf{P}_h}(\mathcal{E}). \quad (16)$$

Using Eqs. (2) and (11) which also stand once the steady state has been achieved, it follows that the mobility is given in the steady state by

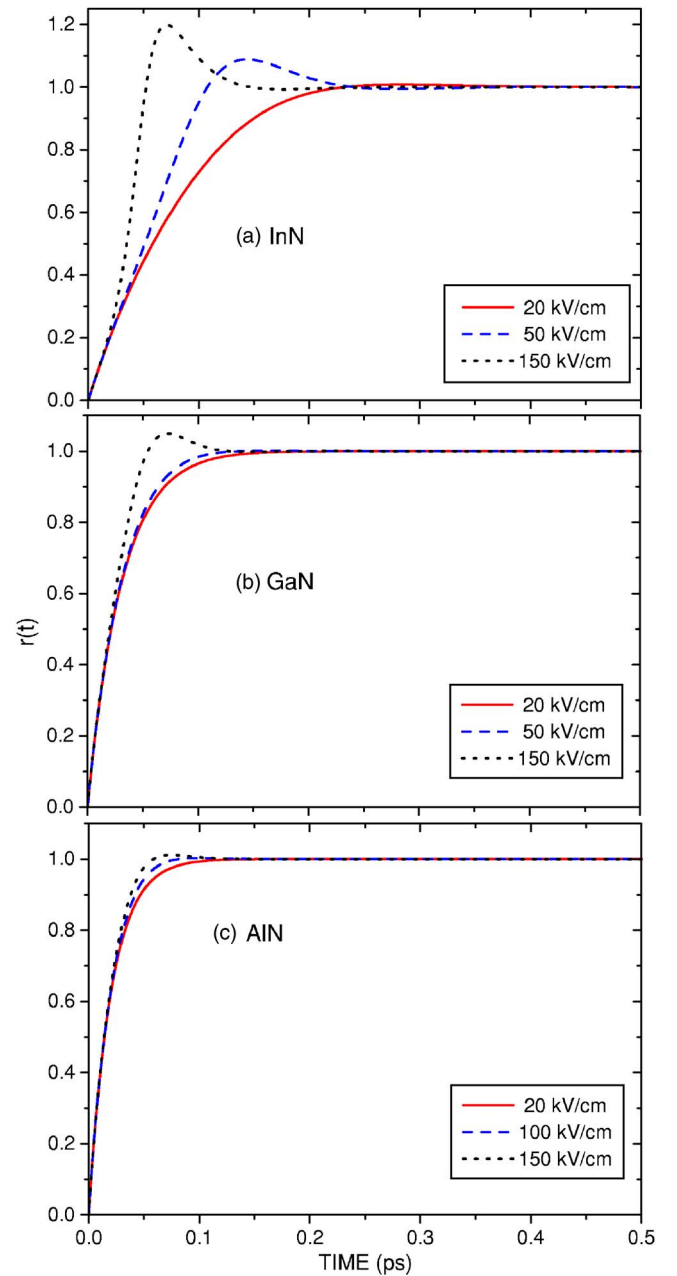


FIG. 1. (Color online) The ratio $r(t) = \tau_{\text{ch}}(t) / \tau_{\mathbf{P}_h}(t)$.

$$\mathcal{M}(\mathcal{E}, n) = \frac{|\mathbf{v}_h(\mathcal{E}, p)|}{\mathcal{E}} = \frac{e}{m_h^*} \tau_{\mathbf{P}_h}(\mathcal{E}, p), \quad (17)$$

where we have stressed that all its dependence on the electric field strength and the concentration is exclusively contained in the momentum relaxation time. As already noticed, such dependence is implicit in the electron and phonon quasitemperatures, $T_h^*(\mathcal{E})$, $T_{\text{Lo}}^*(\mathcal{E})$, and $T_{\text{ac}}^*(\mathcal{E})$, present in $\mathbf{J}_{\mathbf{P}_h}^{(2)}(t)$ and $J_{E_h}^{(2)}(t)$. The detailed expressions for the collision integrals $\mathbf{J}_{\mathbf{P}_h}^{(2)}(t)$ and $J_{E_h}^{(2)}(t)$ are given in Refs. 21 and 37.

III. RESULTS AND DISCUSSION

We solve numerically the set of coupled nonlinear differential Eqs. (1)–(4) to obtain the evolution and steady-state behavior of the basic intensive nonequilibrium thermodynamic variables for wurtzite GaN, AlN, and InN. In these

TABLE I. Parameters of *p*-doped wurtzite AlN, InN, and GaN.

Parameter	AlN	InN	GaN
Hole effective mass m_h^* (m_0)	3.53 ^a	0.5 ^b	2.0 ^a
Lattice parameter a (Å)	3.11 ^c	3.54 ^d	3.19 ^e
Lattice parameter c (Å)	4.98 ^c	5.7 ^d	5.185 ^e
LO phonon energy $\hbar\omega_{LO}$ (meV)	99.2 ^a	89 ^d	92 ^f
Static dielectric constant ϵ_0	8.5 ^c	15.4 ^d	9.5 ^e
Optical dielectric constant ϵ_∞	4.77 ^c	8.4 ^d	5.35 ^e
Long. elast. const. C_l ($\times 10^{12}$ dyn/cm ²)	2.65 ^c	2.65 ^d	2.66 ^g
Trans. elast. const. C_t ($\times 10^{11}$ dyn/cm ²)	4.42 ^c	4.43 ^d	4.41 ^g
Piezoelectric constant h_{pz} (C/m ²)	0.92 ^h	0.375 ^d	0.375 ⁱ

^aReference 39.^bReference 40.^cReference 41.^dReference 42.^eReference 43.^fReference 44.^gReference 45.^hReference 46.ⁱReference 47.

calculations, we consider the contributions to the collision integrals arising out of the different channels of electron scattering, namely, the strong polar-optic (Fröhlich) interaction with Lo phonons, the deformation and piezoelectric interactions with ac phonons, and the interaction with impurities. We have used for these materials the characteristic parameters shown in Table I. The doping concentration is taken as 1.0×10^{18} cm⁻³, and the sample is in contact with a bath at a temperature of 300 K. Numerical results are shown in Figs. 1–8.

The evolution of the hole drift velocity and quasitemperature toward the steady state is depicted in Figs. 2 and 3, respectively. These figures permit to characterize the presence of an overshoot in the hole drift velocity and quasitemperature in InN with the onset of the overshoot effect occurring at ~ 50 kV/cm, which are less pronounced in GaN, with the onset at fields larger than 100 kV/cm, and are practically absent in AlN. We have the fact that at the maximum of the overshoot at, say, time \bar{t} , the time derivative of the hole drift velocity is null and the curvature (second derivative) is negative. Let us consider first the condition of maximum, which implies that at \bar{t} ,

$$\frac{d}{dt} v_h(t)|_{\bar{t}} = \frac{e}{m_h^*} \mathbf{E} - \frac{v_h(\bar{t})}{\tau_{P_h}(\bar{t})} = 0, \quad (18)$$

according to Eqs. (12) and (16). On the other hand, differentiating in time Eq. (13), we have at $t=\bar{t}$ that

$$\frac{d}{dt} v_h(t)|_{\bar{t}} = \frac{e\mathbf{E}}{m_h^*} \frac{d}{dt} \tau_{ch}(t)|_{\bar{t}} = 0; \quad (19)$$

however, using Eq. (14), it follows that

$$\frac{d}{dt} \tau_{ch}(t)|_{\bar{t}} = -\frac{\tau_{ch}(\bar{t})}{\tau_{P_h}(\bar{t})} + 1 = 0, \quad (20)$$

and then, at \bar{t} ,

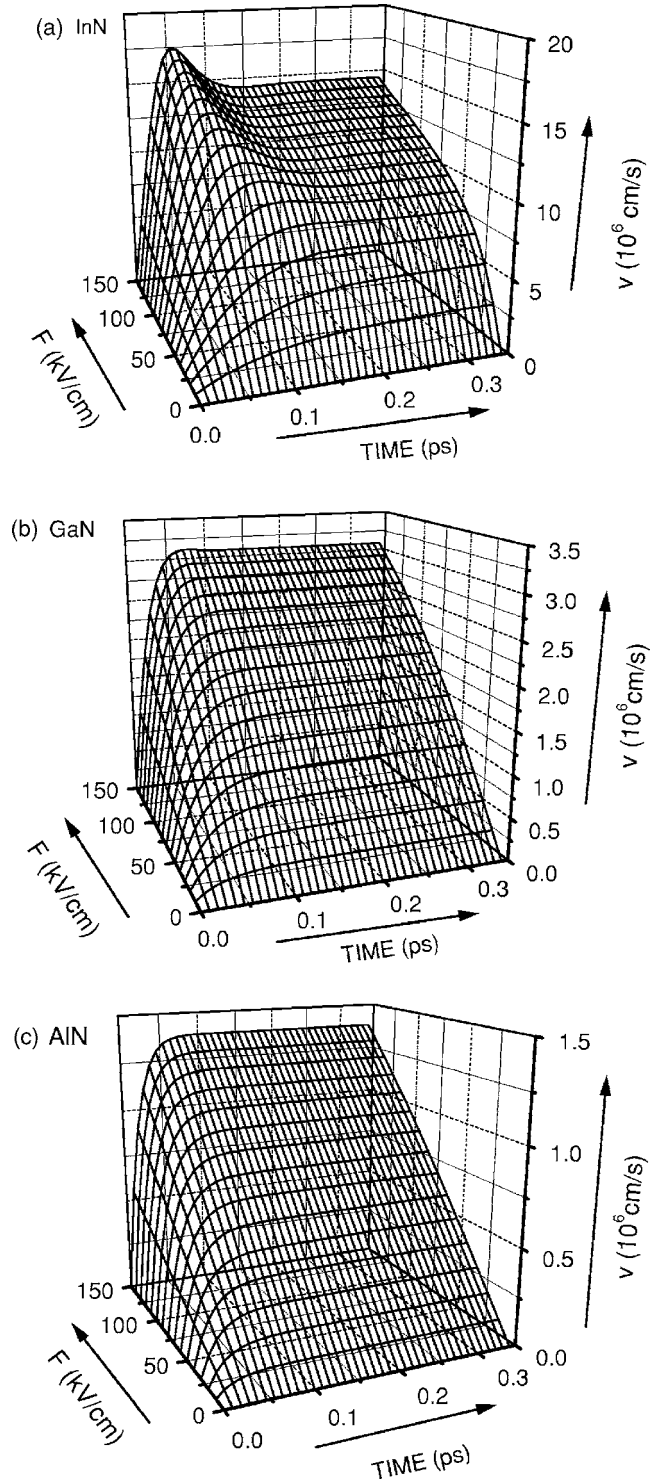


FIG. 2. Time evolution toward the steady state of the hole drift velocity in InN (upper figure), GaN (middle figure), and AlN (lower figure).

$$\tau_{ch}(\bar{t}) = \tau_{P_h}(\bar{t}), \quad (21)$$

that is, the characteristic time for current and the momentum relaxation time are equal, and then the overshoot can only follow in conditions allowing for them to cross (i.e., for sufficiently high values of the electric field strength). This can be seen by inspection of Fig. 1, and also that at low electric field strength, such ratio increases monotonously tending to the value 1 when the steady state is being achieved, and then the momentum relaxation time and the characteristic time for

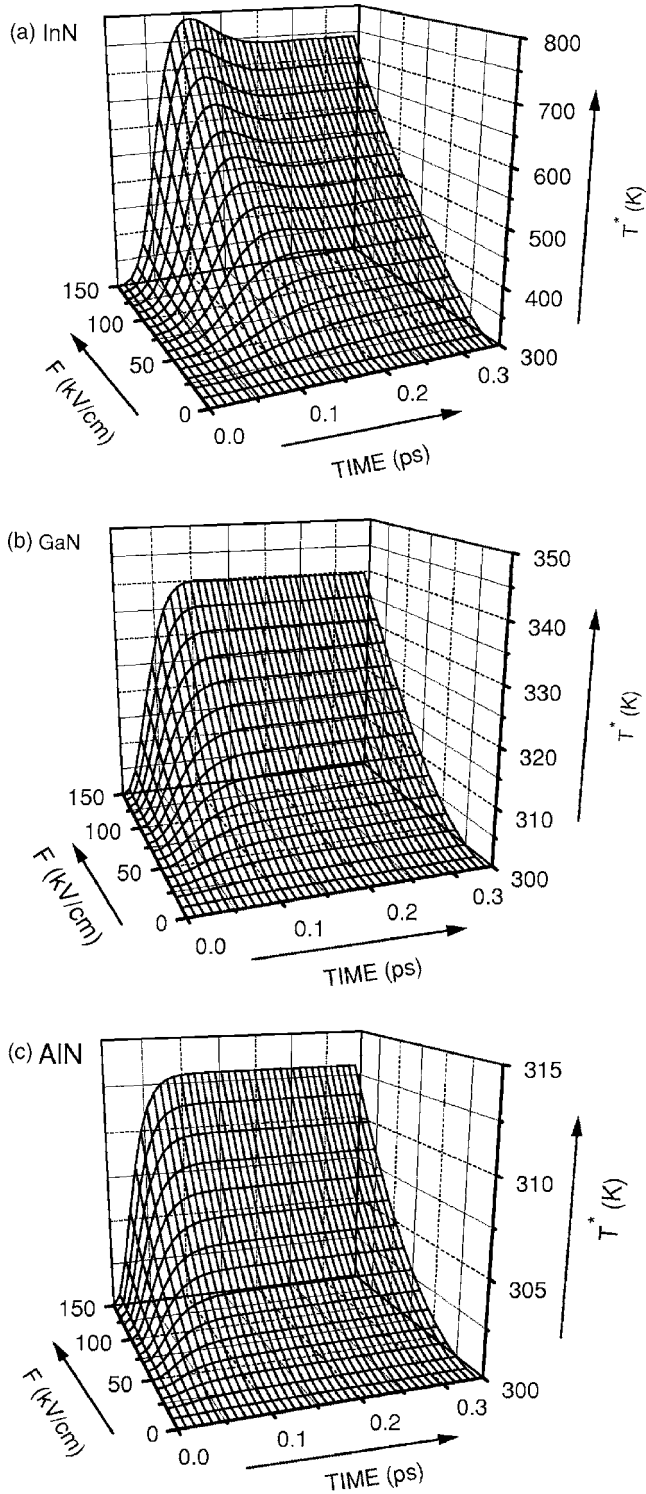


FIG. 3. Time evolution toward the steady state of the hole quasitemperature in InN (upper figure), GaN (middle figure), and AlN (lower figure).

current coincide. It can be noticed that, in the earlier stages of evolution of the nonequilibrium macroscopic (thermodynamic) state of the system, at intermediate to high fields, the characteristic time for current becomes larger than the momentum relaxation time, and there follows velocity overshoot. This is a result of the fact that the momentum relaxation time varies smoothly in time when in the presence of low to intermediate fields, whereas at intermediate to strong fields it changes abruptly decreasing in a subpicosecond time

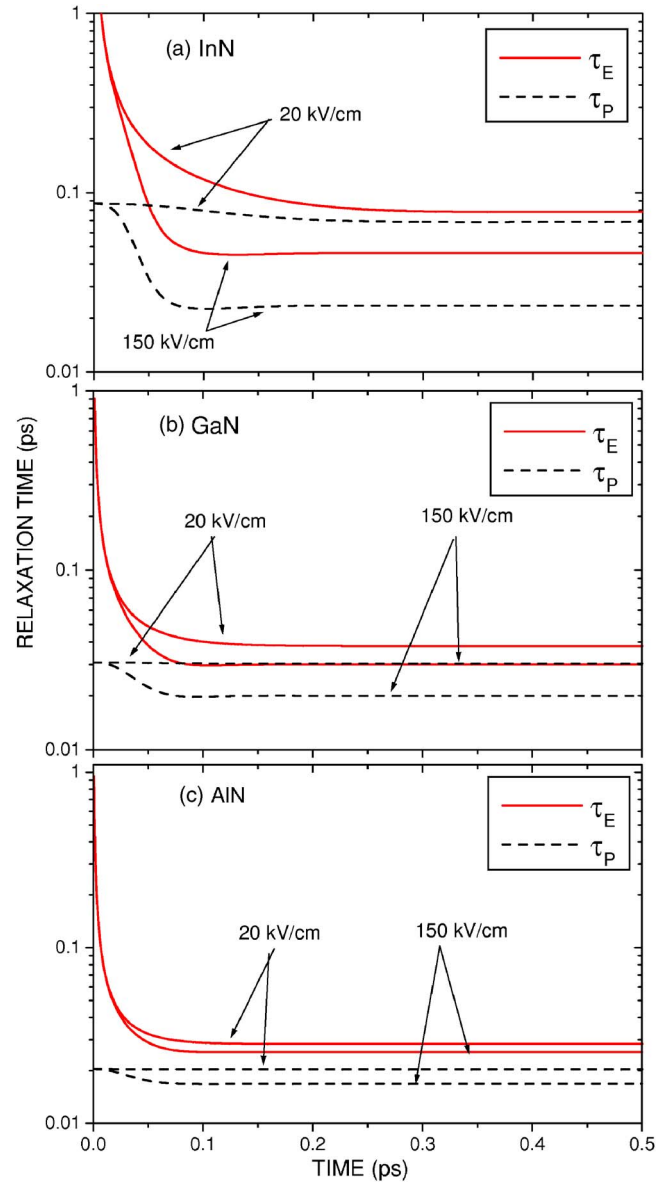


FIG. 4. (Color online) The hole energy relaxation times (solid line), $\tau_{E_h}(t)$, and hole momentum relaxation times (dashed line), $\tau_{P_h}(t)$, for (a) InN, (b) GaN, and (c) AlN.

scale (see Fig. 4). The transition from the, say, normal regime to one displaying overshoot follows around a field of 50 kV/cm for InN.

We recall that the momentum relaxation time is given in Eq. (10), and the energy relaxation time follows from the identification⁴⁸

$$J_{E_h}^{(2)}(t) \equiv - \frac{E_h(t) - E_h^{eq}}{\tau_{E_h}(t)}, \tag{22}$$

where E_h^{eq} is the energy of the holes in equilibrium with the reservoir at temperature T_0 .

The dependence on time of both relaxation times is shown in Fig. 4. The time for the holes to attain the steady state is in our calculations very approximately the same as obtained by other authors, using different descriptions, for the electron in the transient transport regime.¹⁷⁻¹⁹

Figure 5 shows the dependence of the hole drift velocity in the steady state with the electric field intensity. It can be

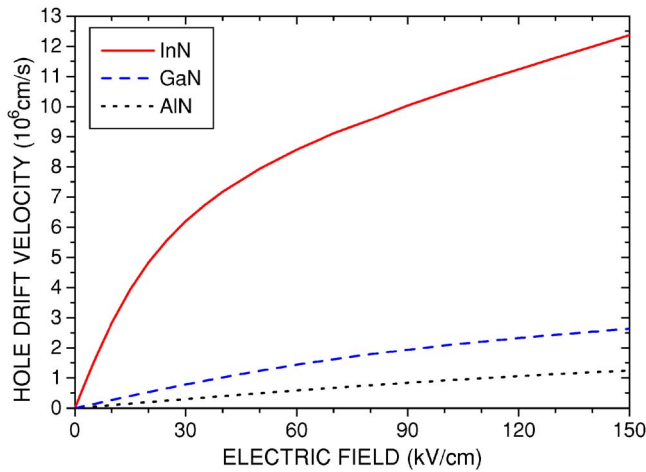


FIG. 5. (Color online) Steady-state hole drift velocity as a function of the electric field in wurtzite p -InN (solid line), p -GaN (dashed line), and p -AlN (dotted line).

seen that at low electric fields, an Ohmic region is present, with a departure from the Ohmic behavior following for larger fields. We can see that InN presents the highest mobility, followed by GaN and AlN. Inspection of Table I tells us that this is a consequence, as expected, of mainly the fact that the effective mass of the holes follows the ordering $m_h^{*(\text{InN})} < m_h^{*(\text{GaN})} < m_h^{*(\text{AlN})}$. We can derive the hole mobility in the steady state, \mathcal{M}_h , as given by $\mathcal{M}_h = |\mathbf{v}_h|/\mathcal{E}$, with the hole drift velocity \mathbf{v}_h related to linear momentum per hole by $\mathbf{P}_h = m_h^* \mathbf{v}_h$. In accordance with Fig. 5, Fig. 6 shows that $\mathcal{M}_h^{(\text{InN})} > \mathcal{M}_h^{(\text{GaN})} > \mathcal{M}_h^{(\text{AlN})}$. It can be noticed that the mobility decreases with increasing strength of the electric field; that is, the differential mobility decreases with increasing field intensity.

Figures 7 and 8 show the dependence with the electric field intensity of the hole and LO-phonon quasitemperatures, respectively, in the steady state. We can verify that the quasitemperature does not increase appreciably above equilibrium temperature in GaN and AlN, however, in InN, the increase is significant (165% in T_h^* and 53% in T_{Lo}^* for an electric field of 150 kV/cm). The quasitemperature of the acoustic phonons, not shown here, has a very small increase, being at most 2% above equilibrium temperature for the

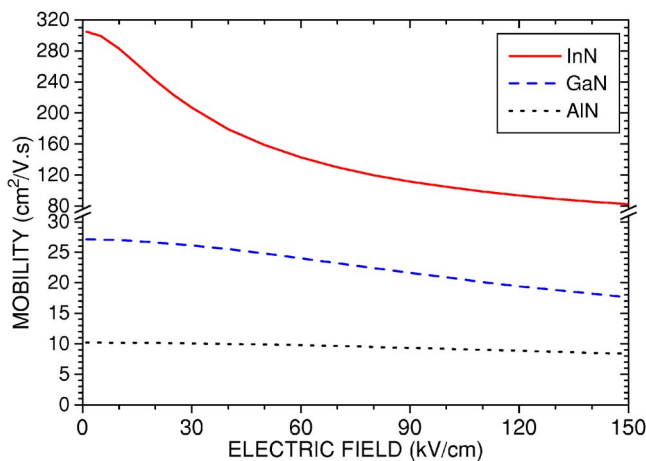


FIG. 6. (Color online) Hole mobility as a function of the electric field in wurtzite p -InN (solid line), p -GaN (dashed line), and p -AlN (dotted line).

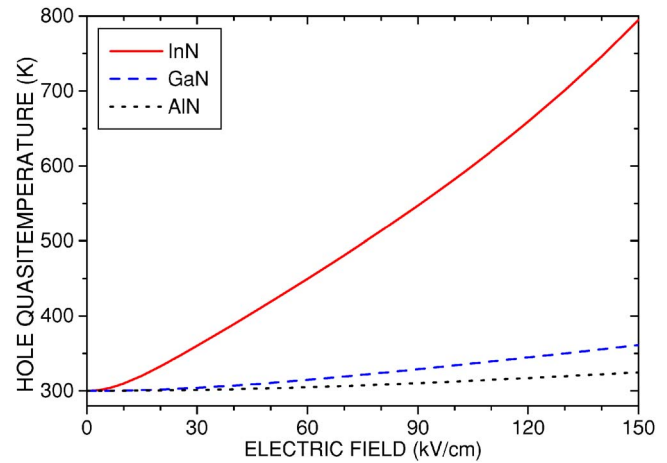


FIG. 7. (Color online) Steady-state hole quasitemperature as function of the electric field in wurtzite p -InN (solid line), p -GaN (dashed line), and p -AlN (dotted line).

higher fields we have considered, a result of the good thermal contact with the reservoir that we have used. Examining Figs. 7 and 8, it can be verified that the values of the hole and LO phonon quasitemperatures increase with a near parabolic law in the electric field intensity, namely, $T^*(F) \approx T_0 + aF + bF^2$, where $T_0 = 300$ K and with the values of the coefficients a and b given in Table II.

IV. SUMMARY AND CONCLUDING REMARKS

In conclusion, we have presented a study on the nonlinear transport in the transient and steady-state regimes of wurtzite p -doped GaN, AlN, and InN; this was done resorting to a quantum kinetic theory derived from the nonequilibrium statistical ensemble formalism. From the point of view of device applications, it is extremely useful to compare the properties of the different materials. The transient transport regime in the III-nitrides follows in a subpicosecond scale (< 0.3 ps), and the existence of an overshoot in the hole drift velocity (and hole quasitemperature) in InN was demonstrated, whose onset occurs at ~ 50 kV/cm. The overshoot effect is a result of the relation between the carrier relaxation rate of momentum and energy: overshoot follows if during

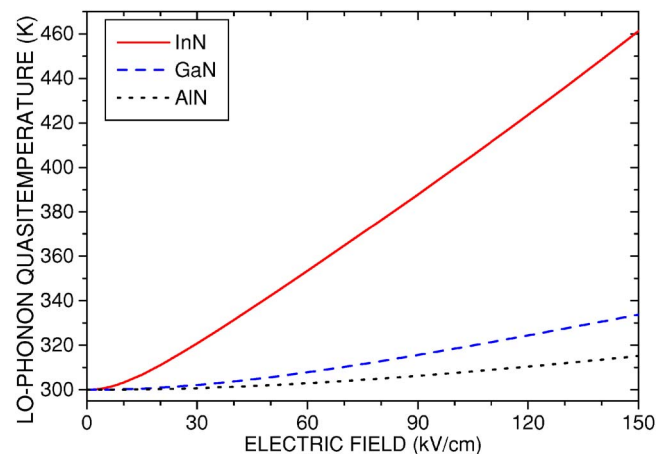


FIG. 8. (Color online) Steady-state LO-phonon quasitemperature as function of the electric field in wurtzite p -InN (solid line), p -GaN (dashed line), and p -AlN (dotted line).

TABLE II. Values of the coefficients a and b .

...	...	a [K(cm/kV)]	b [K(cm/kV) ²]
T_h^*	InN	2.0575	0.0084
	GaN	0.1848	0.0016
	AlN	0.0365	0.0009
T_{LO}^*	InN	0.8586	0.0017
	GaN	0.0931	0.0009
	AlN	0.0205	0.0006

the evolution of the macroscopic state of the system, under the action of the electric field, the former is larger than the latter. The hole mobility was calculated, and it was shown that the larger corresponds to InN, and the smaller to AlN.

It may be noticed that InN shows certain similarities of behavior in its properties with the semiconductor GaAs. The latter was extensively studied some decades ago, particularly in the highly excited condition of a photoinjected plasma of electron and hole pairs. Earlier studies of the ultrafast transient response, and with it an analysis of the relaxation processes developing in such photoinjected plasma, are described in Ref. 48. Ultrafast mobility transient and the steady state that emerges in sequence in the presence of electric fields have been considered (for example, Refs. 49 and 50); and similarly, more recently for III-N.²¹

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