Ultrafast mobility in photoinjected polar semiconductors

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An analytical study of the ultrafast-mobility transient of central-valley nonequilibrium carriers in a highly photoexcited plasma in semiconductors is presented. General expressions for the mobility of the photoinjected carriers are derived. Numerical results are obtained in the case of low to moderately high fields in GaAs. We show that the mobility transient has a structure (maxima and minima) depending on the degree of photoexcitation and electric field intensity. Three different regimes are present, corresponding to (i) structure without overshoot and an Ohmic steady state, (ii) structure with overshoot and a non-Ohmic steady state, and (iii) normal evolution and an Ohmic steady state. A brief discussion of the diffusion coefficient is given.

I. INTRODUCTION

Studies of the optical and transport properties of semiconductors under high levels of excitation have shown novel and quite interesting features, evidenced in ultrafast-laser-spectroscopy experiments. Notable improvements in time-resolved laser spectroscopy have made it a very useful tool to be used with confidence for the investigation of very rapid mechanisms and effects in the biological and physical realms.[1] These kinds of studies are of great interest because of the variety of phenomena observed, most of them of relevance in the functioning of some semiconductor devices, and also because they provide an excellent testing ground for theoretical ideas in the field of many-body systems far from equilibrium. The question of the influence of very fast relaxation processes in highly photoexcited plasma in semiconductors (HEPS) on their optical properties has been the object of experimental and theoretical study. In the present paper we consider how relaxation effects in HEPS affect the ultrafast mobility of the nonequilibrium carriers. Several approaches to the hot-carrier quantum transport are presently available, and numerical methods, such as the Monte Carlo computational approach, have shown recent remarkable improvements.[2] However, analytical methods for studying the nonlinear transport in HEPS under the action of intense electric fields are also desirable to obtain physical insights and for the interpretation of new phenomena. For that purpose we resort here to the use of the nonequilibrium statistical operator method (NSOM). It is a powerful formalism that seems to offer an elegant and concise way for an analytical treatment in the theory of irreversible processes, adequate to deal with a large class of experimental situations. It can be considered a far-reaching generalization of the Chapman-Enskog approach in the kinetic theory of gases[3] or of the Mori-Langevin formalism.[4] The NSOM and its associated generalized nonlinear quantum transport theory are reviewed elsewhere.[4,5] The NSOM we use for the study of nonlinear ultrafast transient transport in HEPS arbitrarily away from equilibrium and for any value of the electric field intensity is that based on Zubarev's approach,[6] in the so-called NSOM linear theory of relaxation (LTR). Numerical calculations appropriate for the case of GaAs are presented. The existence of a novel feature is demonstrated in the evolution curves of the drift velocity, termed *structured ultrafast transport,* viz., maxima—relative or absolute (overshoot)—and minima that may appear before a steady state is reached. A criterion for the occurrence of this structure, and also overshoot effects, is derived.[7] It is shown that there exist three different regimes in the transient transport depending on the range of values of the electric field intensity.

II. NONLINEAR QUANTUM TRANSPORT IN HEPS

We consider the case of a polar semiconductor described by a two-inverted-parabolic-band model, where a concentration n of electron-hole pairs has been created by an intense pulse of laser light. These carriers are in a state strongly departed from equilibrium but in a condition of internal equilibrium (hot carriers) as a result of the Coulomb interaction.[8] A constant electric field of intensity ϵ in, say, the x direction is applied, accelerating the carriers which, at the same time, transfer energy and momentum to the phonon field. The sample is in contact with a thermal reservoir at temperature T_0, and the phonons are warmed up in scattering events involving Fröhlich, deformation potential, and piezoelectric interactions with the carriers.[9]

To deal with the irreversible thermodynamic evolution and transport properties of this system we resort, as indicated in the Introduction, to the NSOM in Zubarev's approach. We recall that the NSOM requires, as a first

step, the choice of a basic set of variables to describe the macrostate of the nonequilibrium system.[4,6,10] For the present case we select the eight dynamical quantities P_j, $j=1,2,\ldots,8$, consisting of the carriers, longitudinal-optical (LO), transverse-optical (TO), and acoustic (AC) phonon Hamiltonians (H_c, $H_{\rm LO}$, $H_{\rm TO}$, and $H_{\rm AC}$, respectively); the number operators for electrons, N_e, and for holes, N_h; and the linear momenta (their components in the direction of the electric field) of electrons, P_e, and of holes, P_h. The nonequilibrium macroscopic variables, i.e., the average values of these eight dynamical quantities over the nonequilibrium ensemble, $Q_j(t)={\rm Tr}[P_j\rho(t)]$, are—after dividing by the volume of the system—the densities of the corresponding energies, $E_c(t)$, $E_{\rm LO}(t)$, $E_{\rm TO}(t)$, $E_{\rm AC}(t)$, the density of pairs, $n(t)$ (equal to the density of the electrons and of the holes), and the density of the linear momenta, $\pi_e(t)$ and $\pi_h(t)$. Further, there are eight intensive nonequilibrium variables $F_j(t)$, thermodynamically conjugated to the $Q_j(t)$, that are defined as the four reciprocal quasitemperatures $\beta_c(t)=1/kT_c^*(t)$, $\beta_{\rm LO}(t)=1/kT_{\rm TO}^*(t)$, $\beta_{\rm TO}(t)=1/kT_{\rm TO}^*(t)$, $\beta_{\rm AC}(t)=1/kT_{\rm AC}^*(t)$; two are associated to quasichemical potentials, $-\beta_c(t)\mu_e(t)$ and $-\beta_c(t)\mu_h(t)$, and, finally, two are associated to drift velocities, $-\beta_c(t)v_e(t)$ and $-\beta_c(t)v_h(t)$. Further analysis of this choice and its limitations is presented in the last section.

The total Hamiltonian H of the HEPS is separated in the form $H=H_0+H'+H_\epsilon$, where H_0 contains the energy operators of each individual subsystem, i.e., $H_0=H_c+H_{\rm LO}+H_{\rm TO}+H_{\rm AC}$, and the interactions between them and with the external reservoirs are included in H', i.e., it is composed of the interaction energies of carriers with the phonon field, anharmonic interaction between phonons, and the interaction with the thermal reservoir (responsible for heat diffusion out of the sample). Finally, the interaction of the carriers with the electric field is

$$H_\epsilon=-e\epsilon\sum_j(x_{ej}-x_{hj})\ , \tag{1}$$

where $x_{e(h)j}$ is the coordinate of the jth electron (hole).

For the given choice of the basic set of dynamical variables we find $[P_j,H_0]=0$ and $[P_j,P_k]=0$, and the Coulomb interaction between carriers, contained in H_c, is treated in the random-phase approximation, i.e., the carriers are considered as a two-component Landau Fermi fluid. Coulomb interaction is only called forth indirectly to ensure the internal thermalization of carriers at any time.[8] The conduction and (heavy-hole) valence Bloch bands are taken in the effective-mass approximation, and we use Einstein models for the optical phonons (with dispersionless frequencies $\omega_{\rm LO}$ and $\omega_{\rm TO}$) and a Debye model for acoustic phonons [$\omega_{\rm AC}(q)=sq$, with triple degeneracy; s is the velocity of sound].

Applying Zubarev's NSOM nonlinear quantum transport theory in the NSOM linear theory of relaxation,[11,4,5,10] we obtain for the evolution equations of the basic set of macrovariables $Q_j(t)$,

$$\frac{1}{V}\frac{d\langle H_C|t\rangle}{dt}=\sum_\alpha\frac{e\epsilon}{m_\alpha v}\langle P_\alpha|t\rangle-\sum_{\alpha,\eta,i}\dot E^i_{\alpha,\eta}(t)\ , \tag{2a}$$

$$\frac{1}{V}\frac{d\langle H_{\rm LO}|t\rangle}{dt}=\sum_{\alpha,i}\dot E^i_{\alpha,{\rm LO}}(t)-\dot E_{\rm LO,AN}(t)\ , \tag{2b}$$

$$\frac{1}{V}\frac{d\langle H_{\rm TO}|t\rangle}{dt}=\sum_{\alpha,i}\dot E^i_{\alpha,{\rm TO}}(t)-\dot E_{\rm TO,AN}(t)\ , \tag{2c}$$

$$\frac{1}{V}\frac{d\langle H_{\rm AC}|t\rangle}{dt}=\sum_{\alpha,i}\dot E^i_{\alpha,{\rm AC}}(t)+\dot E_{\rm LO,AN}(t)+\dot E_{\rm TO,AN}(t)-\dot E_{\rm AC,dif}(t)\ , \tag{2d}$$

$$\frac{1}{V}\frac{d\langle P_e|t\rangle}{dt}=ne\epsilon-\sum_{\eta,i}\dot\pi^i_{e,\eta}(t)\ , \tag{2e}$$

$$\frac{1}{V}\frac{d\langle P_h|t\rangle}{dt}=ne\epsilon-\sum_{\eta,i}\dot\pi^i_{h,\eta}(t)\ , \tag{2f}$$

where we took the modulus of the linear momentum in the direction of the electric field. The two equations for the average number of electrons and holes are not considered since, for our purposes here, they are constant because recombination effects are relevant in a near-nanosecond time scale while our interest is in the time scale of a few picoseconds. In these equations V is the volume of the system, and $\alpha=e$ or h for electrons or holes.

In Eq. (2a) the first term on the right-hand side accounts for the energy transfer from the electric field to the carriers' system, which is the only new term in this equation in comparison with that obtained with a similar derivation in studies of relaxation effects in HEPS.[10] This latter reference also contains an explicit expression for the rate of change of the LO-phonon's energy to the carriers' system. Similar expressions are valid for the two others (TO and AC), namely

$$\dot E^i_{\alpha,\eta}(t)=\frac{2\pi}{\hbar}\sum_{\mathbf{k},\mathbf{q}}\hbar\omega_{\mathbf{q},\eta}|M^i_{\alpha,\eta}(\mathbf{q})|^2\{\nu_{\mathbf{q},n}(t)f_{\mathbf{k},\alpha}(t)[1-f_{\mathbf{k}+\mathbf{q},\alpha}(t)]-[1+\nu_{\mathbf{q},\eta}(t)]f_{\mathbf{k}+\mathbf{q},\alpha}(t)[1-f_{\mathbf{k},\alpha}(t)]\}\delta(\varepsilon_{\mathbf{k}+\mathbf{q},\alpha}-\varepsilon_{\mathbf{k},\alpha}-\hbar\omega_{\mathbf{q},\eta})\ , \tag{3}$$

where the index η is LO, TO, or AC, and the upper index i refers to the different types of interactions, PD, PZ, and FR for deformation potential, piezoelectric, and Fröhlich interactions, respectively. In Eq. (3)

$$\nu_{\mathbf{q},\eta}(t)=1/\{\exp[\beta_\eta(t)\hbar\omega_{\mathbf{q},\eta}]-1\} \tag{4}$$

are the instantaneous distribution functions for η-type phonons with $\omega_{q\eta}$ the corresponding frequency dispersion

relations, and

$$f_{\mathbf{k},\alpha}(t)=[4\pi^3\hbar^3 n/(2\pi m_\alpha)^{3/2}]\beta_c^{3/2}(t) \times\exp\{-\beta_c(t)[\hbar\mathbf{k}-m_\alpha\mathbf{v}_\alpha(t)]^2/2m_\alpha\} \quad (5)$$

are the carrier distribution functions; at the high excitation levels being considered we can use the above instantaneous Maxwell-Boltzmann distribution, which contains a shift term in the exponential due to the presence of the electric field. Further, the right-hand sides of Eqs. (2b)–(2d), besides the contribution of Eq. (3), contain the rates of energy transfer due to anharmonic processes, $\dot{E}_{\eta,\mathrm{AN}}(t)$, and in Eq. (2d) the term $\dot{E}_{\mathrm{AC,dif}}(t)$ due to heat diffusion; all of them are taken in a relaxation-time approximation, viz.,

$$\dot{E}_{\mathrm{LO(TO),AN}}(t)=\sum_{\mathbf{q}}\hbar\omega_{\mathbf{q},\mathrm{LO(TO)}} \times\frac{\nu_{\mathbf{q},\mathrm{LO(TO)}}(t)-\nu_{\mathbf{q},\mathrm{LO(TO)}}(t,\beta_{\mathrm{AC}})}{\tau_{\mathrm{LO(TO)}}} , \quad (6)$$

where

$$\nu_{\mathbf{q},\mathrm{LO(TO)}}(t,\beta_{\mathrm{AC}})=1/\{\exp[\beta_{\mathrm{AC}}(t)\hbar\omega_{\mathbf{q},\mathrm{LO(TO)}}]-1\} , \quad (7)$$

and $\tau_{\mathrm{LO(TO)}}$ is a phenomenological relaxation time, to be evaluated from Raman scattering linewidths.[12] The same term with a change of signal appears in Eq. (12d), and

$$\dot{E}_{\mathrm{AC,dif}}(t)=\sum_{\mathbf{q}}\hbar\omega_{\mathbf{q},\mathrm{AC}}\frac{\nu_{\mathbf{q},\mathrm{AC}}(t)-\nu_{\mathbf{q},\mathrm{AC}}(\beta_B)}{\tau_{\mathrm{AC}}} , \quad (8)$$

where

$$\nu_{\mathbf{q},\mathrm{AC}}(\beta_B)=1/[\exp(\beta_B\hbar\omega_{\mathbf{q},\mathrm{AC}})-1] \quad (9)$$

β_B is the reciprocal temperature of the reservoir, and τ_{AC} is a phenomenological relaxation time which depends on the diffusion coefficient and the dimensions of the surface of the active volume of the crystal.[13]

Finally, in Eqs. (2e) and (2f) the first term on the right-hand side is a drift force due to the action of the electric field. The second term contributes to the rate of variation of the momentum as a result of collisions with phonons, and is given by

$$\dot{\pi}^i_{\alpha,\eta}(t)=\frac{2\pi}{\hbar}\sum_{\mathbf{k},\mathbf{q}}\hbar q_\epsilon|M^i_{\alpha,\eta}(\mathbf{q})|^2\{\nu_{\mathbf{q},\eta}(t)f_{\mathbf{k},\alpha}(t)[1-f_{\mathbf{k}+\mathbf{q},\alpha}(t)] -[1+\nu_{\mathbf{q},\eta}(t)]f_{\mathbf{k}+\mathbf{q},\alpha}(t)[1-f_{\mathbf{k},\alpha}(t)]\}\delta(\varepsilon_{\mathbf{k}+\mathbf{q},\alpha}-\varepsilon_{\mathbf{k},\alpha}-\hbar\omega_{\mathbf{q},\eta}) +\frac{2\pi}{\hbar}\sum_{\mathbf{k},\mathbf{q}}\hbar q_\epsilon|M^i_{\alpha,\eta}(\mathbf{q})|^2\{[\nu_{\mathbf{q},\eta}(t)+1]f_{\mathbf{k},\alpha}(t)[1-f_{\mathbf{k}-\mathbf{q},\alpha}(t)] -\nu_{\mathbf{q},\eta}(t)f_{\mathbf{k}-\mathbf{q},\alpha}(t)[1-f_{\mathbf{k},\alpha}(t)]\}\delta(\varepsilon_{\mathbf{k}-\mathbf{q},\alpha}-\varepsilon_{\mathbf{k},\alpha}+\hbar\omega_{\mathbf{q},\eta}) , \quad (10)$$

where q_ϵ is the modulus of the component of $\mathbf{q}$ in the direction of the electric field.

Finally, performing in the complete set of equations of evolution the integrations in reciprocal space (summation over $\mathbf{k}$ and $\mathbf{q}$) involved in all the collision operators, we obtain

$$\dot{E}^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)=A^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)\exp[-x_\alpha(t)]\{M(3,\tfrac{3}{2},x_\alpha(t))+[3y_\alpha(t)-\beta_C(t)/\beta_{\mathrm{AC}}(t)]M(2,\tfrac{3}{2},x_\alpha(t)) -(3\pi^{1/2}/2)y_\alpha^{1/2}(t)M(\tfrac{5}{2},\tfrac{3}{2},x_\alpha(t))-\pi^{1/2}y_\alpha^{3/2}(t)\exp[x_\alpha(t)] -\tfrac{1}{3}y_\alpha(t)[1-3\sqrt{2}y_\alpha(t)]M(1,\tfrac{3}{2},x_\alpha(t))\} , \quad (11)$$

$$\dot{E}^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)=A^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)\exp[-x_\alpha(t)]\{M(2,\tfrac{3}{2},x_\alpha(t))+[y_\alpha(t)-\beta_c(t)/\beta_{\mathrm{AC}}(t)]M(1,\tfrac{3}{2},x_\alpha(t))-\pi^{1/2}y_\alpha^{1/2}(t)\exp[x_\alpha(t)]\} , \quad (12)$$

where

$$A^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)=\theta_{\mathrm{AC}}n\left[\frac{2^{7/2}m_\alpha^{5/2}E_{1\alpha}^2}{\pi^{3/2}\hbar^4\rho}\right]\beta_C^{-3/2}(t) , \quad (13)$$

$$A^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)=\theta_{\mathrm{AC}}n\left[\frac{2^{1/2}m_\alpha^{3/2}e^2H_{\mathrm{PZ}}^2}{\pi^{3/2}\hbar^2\rho\epsilon_0^2}\right]\beta_C^{-1/2}(t) , \quad (14)$$

$$x_\alpha(t)=\beta_C(t)\tfrac{1}{2}m_\alpha v_\alpha^2(t) , \quad (15a)$$

$$y_\alpha(t)=\beta_C(t)\tfrac{1}{2}m_\alpha s^2 . \quad (15b)$$

$E_{1\alpha}$ is the deformation potential coupling constant, H_{PZ} the piezoelectric potential coupling constant, ϵ_0 the static dielectric constant, ρ the density of the material, $\theta_{\mathrm{AC}}=3$ is the degeneracy of the acoustic modes, and $M(a,b,x)$ are

Kummer functions;[14] also

$$\dot{E}^{\mathrm{FR}}_{\alpha,\mathrm{LO}}(t)= A^{\mathrm{FR}}_{\alpha,\mathrm{LO}}(t)\exp[-z_{\mathrm{LO}}(t)-x_\alpha(t)] \times\sum_{l=1}^{\infty}\sum_{l'=0}^{\infty}\frac{1}{(2l'+1)!}\frac{[2^3 z_{\mathrm{LO}}(t)x_\alpha(t)]^{l'}}{2l-1} \times\{[1+\nu_{\mathrm{LO}}(t)]\exp[-z_{\mathrm{LO}}(t)]\Gamma(l+\tfrac{1}{2})U(l+\tfrac{1}{2},l'+2;2z_{\mathrm{LO}}(t)) -\nu_{\mathrm{LO}}(t)\exp[z_{\mathrm{LO}}(t)]\Gamma(l'+l+\tfrac{1}{2})U(l'+l+\tfrac{1}{2},l'+2;2z_{\mathrm{LO}}(t))\}\ , \tag{16}$$

$$\dot{E}^{\mathrm{PD}}_{\alpha,\mathrm{LO}}(t)= A^{\mathrm{PD}}_{\alpha,\mathrm{LO}}(t)\exp[-z_{\mathrm{LO}}(t)-x_\alpha(t)] \times\sum_{l=0}^{\infty}\frac{[2^3 z_{\mathrm{LO}}(t)x_\alpha(t)]^{l}}{(2l+1)!}\{[1+\nu_{\mathrm{LO}}(t)]\exp[-z_{\mathrm{LO}}(t)]\Gamma(\tfrac{3}{2})U(\tfrac{3}{2},l+3;2z_{\mathrm{LO}}(t)) -\nu_{\mathrm{LO}}(t)\exp[z_{\mathrm{LO}}(t)]\Gamma(l+\tfrac{3}{2})U(l+\tfrac{3}{2},l+3;2z_{\mathrm{LO}}(t))\}\ , \tag{17}$$

where

$$\nu_{\mathrm{LO}}(t)=1/\{\exp[\beta_{\mathrm{LO}}(t)\hbar\omega_{\mathrm{LO}}]-1\}\ , \tag{18}$$

$$A^{\mathrm{FR}}_{\alpha,\mathrm{LO}}(t)=\theta_{\mathrm{LO}}n\left[\frac{2^{3/2}eE_{0\alpha}}{\pi^{1/2}m_\alpha^{1/2}}\right](\hbar\omega_{\mathrm{LO}})^2\beta_C^{3/2}(t)\ , \tag{19}$$

$$A^{\mathrm{PD}}_{\alpha,\mathrm{LO}}(t)=\theta_{\mathrm{LO}}n\left[\frac{m_\alpha^{3/2}D^2_{\mathrm{LO}\alpha}}{\pi^{3/2}\hbar_\rho^2}\right](\hbar\omega_{\mathrm{LO}})^2\beta_C^{3/2}(t)\ , \tag{20}$$

$$z_{\mathrm{LO}}(t)=\beta_C(t)\hbar\omega_{\mathrm{LO}}/2\ . \tag{21}$$

$E_{0\alpha}$ is the Fröhlich field, the $D_{\mathrm{LO}\alpha}$ are the LO-phonon deformation potential interaction coupling constants, $\theta_{\mathrm{LO}}=1$ is the degeneracy of LO phonons, $\Gamma(a)$ are gamma functions, $U(a,b,z)$ are confluent hypergeometric functions,[14]

$$\dot{E}^{\mathrm{PD}}_{\alpha,\mathrm{TO}},\ \text{the same as}\ \dot{E}^{\mathrm{PD}}_{\alpha,\mathrm{LO}}\ \text{with the exchange LO}\rightarrow\text{TO}\ , \tag{22}$$

and $\theta_{\mathrm{TO}}=2$.

The different contributions to the equations of evolution for the linear momentum are

$$\dot{\pi}^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)=B^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)v_\alpha(t)\exp[-x_\alpha(t)]\{[1-6y_{\mathrm{AC},\alpha}(t)]M(3,\tfrac{5}{2},x_\alpha(t)) +2\pi^{1/2}y^{1/2}_{\mathrm{AC},\alpha}(t)[\beta_{\mathrm{AC}}(t)/\beta_C(t)]^{1/2}M(\tfrac{7}{2},\tfrac{5}{2},x_\alpha(t))+\cdots\}\ , \tag{23}$$

$$\dot{\pi}^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)=B^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)v_\alpha(t)\exp[-x_C(t)]\{[1-2y_{\mathrm{AC},\alpha}(t)]M(2,5/2;x_\alpha(t))+y^{1/2}_{\mathrm{AC},\alpha}(t)[\beta_{\mathrm{AC}}(t)/\beta_C(t)]\exp[x_C(t)]\}\ , \tag{24}$$

where

$$B^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)=\theta_{\mathrm{AC}}n\left[\frac{2^{3/2}m_\alpha^{5/2}E^2_{1\alpha}}{\pi^{3/2}\hbar^4 s^2\rho}\right]\frac{\beta_C^{-1/2}(t)}{\beta_{\mathrm{AC}}(t)}\ , \tag{25}$$

$$B^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)=\theta_{\mathrm{AC}}n\left[\frac{2^{1/2}m_\alpha^{3/2}e^2H^2_{\mathrm{PZ}}}{3\pi\hbar^2 s^2\epsilon_0^2\rho}\right]\frac{\beta^{1/2}(t)}{\beta_{\mathrm{AC}}(t)}\ , \tag{26}$$

$$y_{\mathrm{AC},\alpha}(t)=\beta_{\mathrm{AC}}(t)\tfrac{1}{2}m_\alpha s^2\ . \tag{27}$$

Also,

$$\dot{\pi}^{\mathrm{FR}}_{\alpha,\mathrm{LO}}(t)= B^{\mathrm{FR}}_{\alpha,\mathrm{LO}}(t)x_\alpha^{-1/2}(t)\exp[-2z_{\mathrm{LO}}(t)-x_\alpha(t)] \times\sum_{l=0}^{\infty}\left[\left[\frac{2^{3l}}{(2l+1)!}-\frac{2^{3l}}{(2l)!}\right][x_\alpha(t)z_{\mathrm{LO}}(t)]^l \times\left[\Gamma(l+\tfrac{1}{2})\{[1+\nu_{\mathrm{LO}}(t)]U(\tfrac{3}{2},l+2;2z_{\mathrm{LO}}(t))+\nu_{\mathrm{LO}}(t)U(l+\tfrac{1}{2},l+2;2z_{\mathrm{LO}}(t))\} +\sum_{l'=1}^{\infty}\frac{1}{(2l'-1)!}\{[1+\nu_{\mathrm{LO}}(t)]\Gamma(l'+\tfrac{1}{2})U(l'+\tfrac{1}{2},l+1;2z_{\mathrm{LO}}(t)) -\nu_{\mathrm{LO}}(t)\Gamma(l+l'-\tfrac{1}{2})U(l+l'-\tfrac{1}{2},l+1;2z_{\mathrm{LO}}(t))\}\right]\right]\ , \tag{28}$$

$$\dot{\pi}^{\mathrm{PD}}_{\alpha,\mathrm{LO}}(t)=B^{\mathrm{PD}}_{\alpha,\mathrm{LO}}(t)x_\alpha^{-1/2}(t)\exp[-x_\alpha(t)]$$
$$\times\left\{\left[\left[1+\frac{\nu(t)}{\nu_{\mathrm{LO}}(t)}\right]+\exp[-2z_{\mathrm{LO}}(t)]\left[1-\frac{\nu(t)}{\nu_{\mathrm{LO}}(t)}\right]\right]\right.$$
$$\times\left[\sum_{l=0}^{\infty}\left[\frac{2^{3l}}{(2l+1)!}-\frac{2^{3l}}{(2l)!}\right][x_\alpha(t)z_{\mathrm{LO}}(t)]^l\frac{6z_{\mathrm{LO}}(t)}{2l+1}U(\tfrac{5}{2},l+4;2z_{\mathrm{LO}}(t))\right]$$
$$+\left[\left[1-\frac{\nu(t)}{\nu_{\mathrm{LO}}(t)}\right]+\exp[-2z_{\mathrm{LO}}(t)]\left[1-\frac{\nu(t)}{\nu_{\mathrm{LO}}(t)}\right]\right]$$
$$\times\left\{\sum_{l=0}^{\infty}\left[\frac{2^{3l}}{(2l+1)!}-\frac{2^{3l}}{(2l)!}\right][x_\alpha(t)z_{\mathrm{LO}}(t)]^l\right.$$
$$\left.\left.\times\left[U(\tfrac{3}{2},l+2;2z_{\mathrm{LO}}(t))+\left[1-\frac{3}{2l+1}\right]U(\tfrac{3}{2},l+3;2z_{\mathrm{LO}}(t))\right]\right\}\right\}, \tag{29}$$

where

$$B^{\mathrm{FR}}_{\alpha,\mathrm{LO}}(t)=\theta_{\mathrm{LO}}n\left[\frac{eE_{0\alpha}}{\pi^{1/2}}\right]\hbar\omega_{\mathrm{LO}}\beta_C(t)\ , \tag{30}$$

$$B^{\mathrm{PD}}_{\alpha,\mathrm{LO}}(t)=\theta_{\mathrm{LO}}n\left[\frac{m_\alpha^2E_{10\alpha}^2}{2^2\pi\hbar^4s^2\rho}\right]\times(\hbar\omega_{\mathrm{LO}})^{5/2}\nu_{\mathrm{LO}}(t)\beta_C(t)\ , \tag{31}$$

$$\nu(t)=1/\{\exp[\beta_C(t)\hbar\omega_{\mathrm{LO}}]-1\}\ , \tag{32}$$

and finally,

$\dot{\pi}^{\mathrm{FR}}_{\alpha,\mathrm{TO}}$, the same as $\dot{\pi}^{\mathrm{FR}}_{\alpha,\mathrm{LO}}$ with the exchange LO↔TO , (33)

$\dot{\pi}^{\mathrm{PD}}_{\alpha,\mathrm{TO}}$, the same as $\dot{\pi}^{\mathrm{PD}}_{\alpha,\mathrm{LO}}$ with the exchange LO↔TO . (34)

So far we have obtained a complete derivation of the equations that govern the evolution of the basic set of variables that describe the macrostate of the photoinjected HEPS in a constant electric field. They are valid for any intensity of the electric field strength but relaxation effects due to collision with phonons have been treated in the NSOM linear theory of relaxation. The collision operators are expressed in terms of series of Kummer and confluent hypergeometric functions, and thus are rather difficult to manipulate numerically. To simplify these expressions we restrict the calculations to the case of low to moderately high electric fields assuming that the carrier energy of drift is smaller than or, at most, comparable to the thermal energy, i.e.,

$$\beta_c(t)m_\alpha v_\alpha^2(t)/2\lesssim 1\ . \tag{35}$$

In these conditions, the series in Eqs. (11), (2), (16), (17), and (22) can be rearranged in the form of a dominant term plus corrections. We retain only the main terms to obtain

$$\dot{E}^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)=A^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)\left[1-\frac{\beta_C(t)}{\beta_{\mathrm{AC}}(t)}\right], \tag{36}$$

$$\dot{E}^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)=A^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)\left[1-\frac{\beta_C(t)}{\beta_{\mathrm{AC}}(t)}\right], \tag{37}$$

$$\dot{E}^{\mathrm{FR}}_{\alpha,\mathrm{LO}}(t)=n\theta_{\mathrm{LO}}eE_{0\alpha}\left[\frac{4\hbar\omega_{\mathrm{LO}}}{\pi m_\alpha}\right]^{1/2}\left[1-\frac{\nu_{\mathrm{LO}}(t)}{\nu(t)}\right]\times z_{\mathrm{LO}}^{1/2}(t)\exp[-z_{\mathrm{LO}}(t)]K_0(z_{\mathrm{LO}}(t))\ , \tag{38}$$

$$\dot{E}^{\mathrm{PD}}_{\alpha,\mathrm{LO}}(t)=n\theta_{\mathrm{LO}}D^2_{\mathrm{LO},\alpha}\left[\frac{2m_\alpha^3\hbar\omega_{\mathrm{LO}}}{\pi^3\hbar^4\rho^2}\right]\left[1-\frac{\nu_{\mathrm{LO}}(t)}{\nu(t)}\right]\times z_{\mathrm{LO}}^{1/2}(t)\exp[-z_{\mathrm{LO}}(t)]K_1(z_{\mathrm{LO}}(t))\ , \tag{39}$$

$\dot{E}^{\mathrm{PD}}_{\alpha,\mathrm{TO}}(t)$, the same as $\dot{E}^{\mathrm{PD}}_{\alpha,\mathrm{LO}}(t)$ with the exchange LO↔TO (40)

for the terms associated with relaxation of energy (we used also that $\beta_{\mathrm{AC}}hsq\ll 1$), and for the terms associated with momentum relaxation we find

$$\dot{\pi}^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)=B^{\mathrm{PD}}_{\alpha,\mathrm{AC}}(t)v_\alpha(t)\ , \tag{41}$$

$$\dot{\pi}^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)=B^{\mathrm{PZ}}_{\alpha,\mathrm{AC}}(t)v_\alpha(t)\ , \tag{42}$$

$$\dot{\pi}^{\mathrm{FR}}_{\alpha,\mathrm{LO}}(t)=\theta_{\mathrm{LO}}nm_\alpha\gamma_\alpha(t)v_\alpha(t)\ , \tag{43}$$

$$\dot{\pi}^{\rm PD}_{\alpha,{\rm LO}}(t)=\theta_{\rm LO}n\left[\frac{2^{-1/2}m_\alpha^{3/2}E_{10\alpha}^2}{3\pi^{3/2}\hbar^4 s^2\rho}\right](\hbar\omega_{\rm LO})^3\nu_{\rm LO}(t)\beta_C^{3/2}(t)\exp[z_{\rm LO}(t)]v_\alpha(t)$$
$$\times\left\{\left[\left[1+\frac{\nu(t)}{\nu_{\rm LO}(t)}\right]+\exp[-2z_{\rm LO}(t)]\left[1-\frac{\nu(t)}{\nu_{\rm LO}(t)}\right]\right]K_2(z_{\rm LO}(t))\right.$$
$$\left.-\left[\left[1-\frac{\nu(t)}{\nu_{\rm LO}(t)}\right]-\exp[-2z_{\rm LO}(t)]\left[1-\frac{\nu(t)}{\nu_{\rm LO}(t)}\right]\right]K_1(z_{\rm LO}(t))\right\}, \tag{44}$$

$$\dot{\pi}^{\rm PD}_{\alpha,{\rm TO}}(t)=\text{same as }\dot{\pi}^{\rm PD}_{\alpha,{\rm LO}}(t)\text{ exchanging LO}\leftrightarrow\text{TO}\ , \tag{45}$$

where

$$\gamma_\alpha(t)=\gamma_{\alpha 0}x_\alpha^{3/2}(t)\exp[z_{\rm LO}(t)]\nu_{\rm LO}(t)\left\{\left[\left[1+\frac{\nu(t)}{\nu_{\rm LO}(t)}\right]+\exp[-2z_{\rm LO}(t)]\left[1-\frac{\nu(t)}{\nu_{\rm LO}(t)}\right]\right]K_1(z_{\rm LO}(t))\right.$$
$$\left.-\left[\left[1-\frac{\nu(t)}{\nu_{\rm LO}(t)}\right]-\exp[-2z_{\rm LO}(t)]\left[1-\frac{\nu(t)}{\nu_{\rm LO}(t)}\right]\right]K_0(z_{\rm LO}(t))\right\} \tag{46}$$

with

$$\nu(t)=1/\{\exp[\beta_C(t)\hbar\omega_{\rm LO}]-1\}\ , \tag{47}$$

$$\gamma_{\alpha 0}=(2^{5/2}eE_{0\alpha}/3)(1/2\pi m_\alpha\hbar\omega_{\rm LO})^{1/2}\ , \tag{48}$$

and $K_n(z)$ are Bessel functions of second order.[14]

Simple mathematical manipulations allow us to put Eq. (44) in the form given by Conwell[9] with the instantaneous values of $T_c(t)$ and $T_{\rm LO}(t)$.

Next, we assume that the Fröhlich interaction predominates over all other carrier-phonon interactions, and, since in the very early stages of relaxation after finalization of the laser pulse there is practically no heating of A phonons, we take $T_{\rm AC}$ equal to the reservoir temperature. Hence, the original set of six generalized transport equations reduce to four equations once those for the AC and TO phonons' rate of energy variation are dropped.

The left-hand sides of Eqs. (2) are expressed in terms of the intensive variables, i.e., quasitemperatures and drift velocities, using the relations

$$E_c(t)=\sum_{\mathbf{k}\alpha}\epsilon_{\mathbf{k},\alpha}f_{\mathbf{k},\alpha}(t)\ , \tag{49a}$$

$$E_\eta(t)=\sum_{\mathbf{q}}\theta_\eta\hbar\omega_{\mathbf{q},\eta}\nu_{\mathbf{q},\eta}(t)\ , \tag{49b}$$

$$\frac{1}{V}\langle P_\alpha|t\rangle=nm_\alpha v_\alpha(t)\ , \tag{49c}$$

and then

$$\frac{d\beta_C(t)}{dt}=-\frac{\beta_C^2(t)}{3}\left[\sum_\alpha\frac{v_\alpha(t)}{n}\dot{\pi}^{\rm FR}_{\alpha,{\rm LO}}(t)-\dot{E}^{\rm PD}_{\alpha,{\rm AC}}(t)-\dot{E}^{\rm FR}_{\alpha,{\rm LO}}(t)\right], \tag{50a}$$

$$\frac{d\beta_{\rm LO}(t)}{dt}=\frac{2V_{\rm cell}}{(\hbar\omega_{\rm LO})^2}\{1-\cosh[\beta_{\rm LO}(t)\hbar\omega_{\rm LO}]\}\times\left[\sum_\alpha\dot{E}^{\rm FR}_{\alpha,{\rm LO}}(t)-\dot{E}_{\rm LO,AN}(t)\right], \tag{50b}$$

$$\frac{dv_\alpha(t)}{dt}=(e/m_\alpha)\epsilon-\gamma_\alpha(t)v_\alpha(t)\ . \tag{50c}$$

Equation (50c) is a Newton-Langevin-type equation with $\gamma_\alpha(t)$ playing the role of the reciprocal of an instantaneous momentum relaxation time. Equations (50c) and (50d) are of the same form as Eqs. (42) in Ref. 5, but γ_α in the latter depends on a supercorrelation function [Eq. (43) in Ref. 5], a functional with a highly complicated dependence of all the nonequilibrium variables, including the drift velocities. Differently, $\gamma_\alpha(t)$ of Eqs. (50c) and (50d), calculated in the NSOM linear theory of relaxation, is dependent only on $T_c(t)$ and $T_{\rm LO}(t)$, but is independent of $v_\alpha(t)$. Thus, Eqs. (50c) and (50d) are first-order linear differential equations for each drift velocity, possessing the solutions

$$v_\alpha(t)=(e/m_\alpha)\epsilon\tau_\alpha(t)\ , \tag{51}$$

where

$$\tau_\alpha(t)=e^{-\psi(t)}\int_0^t dt'e^{\psi(t')}\ , \tag{52a}$$

$$\psi_\alpha(t)=\int_0^t dt'\gamma_\alpha(t')\ , \tag{52b}$$

and we have taken the initial condition $v_\alpha(0)=0$.

Defining the currents $I_\alpha(t)=nev_\alpha(t)$ and using Eq. (51) we obtain a Drude-type conductivity

$$\sigma_\alpha(t)=(ne^2/m_\alpha)\tau_\alpha(t)\ , \tag{53}$$

with an instantaneous transport relaxation time depending on time through the quasitemperatures $T_c(t)$ and $T_{\rm LO}(t)$, and then varying in time with the irreversible evolution of the macrostate of the system.

Using Eqs. (51) and (52) straightforwardly proves that the drift velocity would have extremal points at, say, a time t_x where

$$\left.\frac{dv_\alpha}{dt}\right|_{t_x}=(e/m_\alpha)\epsilon[1-\gamma(t_x)\tau(t_x)]=0\ , \tag{54}$$

i.e., whenever there occurs during the transient period a

crossover of the evolution curves for the momentum relaxation time and transport relaxation time. Such an extremum is a maximum or a minimum if the second time derivative at t_x,

$$\left.\frac{d^2v_\alpha}{dt^2}\right|_{t_x} = -(e/m_\alpha)\epsilon\tau_\alpha(t_x)\left.\frac{d\gamma_\alpha}{dt}\right|_{t_x}$$
$$= v_\alpha(t_x)\gamma_\alpha^2(t_x)\left.\frac{d\gamma_\alpha^{-1}}{dt}\right|_{t_x}, \tag{55}$$

is negative or positive, respectively. Further,

$$\left.\frac{d\gamma_\alpha^{-1}}{dt}\right|_{t_x} = \frac{\partial\gamma_\alpha^{-1}}{\partial T_c}\left.\frac{dT_c}{dt}\right|_{T_x} + \frac{\partial\gamma^{-1}}{\partial T_{\rm LO}}\left.\frac{dT_{\rm LO}}{dt}\right|_{t_x} \tag{56}$$

and, since the last term is expected to be much smaller than the first,[10] neglecting it we find that on cooling ($dT_c/dt<0$) a maximum occurs if $\partial\gamma^{-1}/T_c>0$, and a minimum for $\partial\gamma^{-1}/T_c<0$. Once $\tau(t)$ begins at zero and increases, the first extremum, if it occurs, is a maximum, and thus a *transient with structure (a maximum and a minimum) should follow if on cooling γ^{-1} passes through a minimum.* It ought to be emphasized that this characteristic of the mobility transient remains valid for the quite general case of any intensity of the electric field strength and large relaxation effects, as shown in Ref. 5. We recall that the expression for $\gamma(t)$ given by Eq. (46) is valid for low to moderately high fields and the NSOM linear theory of relaxation. Also, we call attention to the fact that the maximum is an overshoot if the momentum relaxation time at t_x is larger than γ^{-1} at the stationary state.

We have drawn in Fig. 1 the curves showing the dependence of the reciprocal of the momentum relaxation time on the carriers' quasi-temperature for several values of the LO-phonons' quasitemperature. In it $\Theta_0=\hbar\omega_0/k$ is the Einstein temperature, and γ_α is normalized in terms of the $\gamma_{\alpha 0}$ of Eq. (48). The existence of a maximum of γ can be seen (minimum of the momentum relaxation time γ^{-1}) for a quasitemperature $T_c^{\rm ex}\simeq 2\Theta_0$, with a negative slope (or alternatively $d\gamma^{-1}/dt_c>0$) for $T_c>T_c^{\rm ex}$ and a positive slope (or $d\gamma^{-1}/dT_c<0$) for $T_c<T_c^{\rm ex}$. It can also be seen that $d\gamma/dT_{\rm LO}$ is smooth, and thus the condition for the possible occurrence of a structured transient is fulfilled.

Next we apply these results to a specific case to obtain numerical solutions.

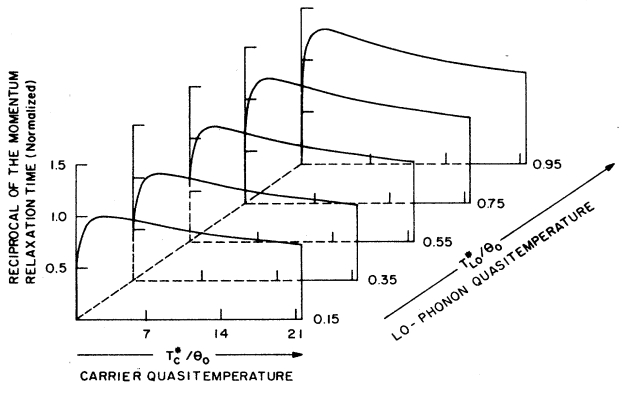

FIG. 1. Dependence of the reciprocal of the momentum relaxation time with the carrier quasitemperature for several values of the LO-phonon quasitemperature. The normalization factors are indicated in the main text.

III. STRUCTURED MOBILITY IN GaAs

Consider a sample of GaAs illuminated by an intense pulse of laser light. To fix initial conditions, we take as an example the case of the experiment of Shank *et al.*[15] which we previously used to study relaxation phenomena in HEPS.[10] A very short laser pulse of 0.25 psec produces a density of photoinjected carriers $n=2\times10^{18}$ cm^{-3}, having an excess kinetic energy of roughly 2.4 eV, and in contact with a thermal reservoir at 300 K. Immediately after the pulse, the initial quasitemperature of carriers is, roughly, 6700 K, and $T_{\rm LO}\sim303$ K, $T_{\rm TO}\sim303$ K, $T_{\rm AC}\sim T_0=300$ K.[10] We have also used relaxation times for anharmonic processes of 10 psec and a heat diffusion relaxation time of 1 nsec.

The coupled set of differential equations, Eqs. (50), is solved using standard computational techniques. Our results are displayed in the accompanying figures. Figure 2

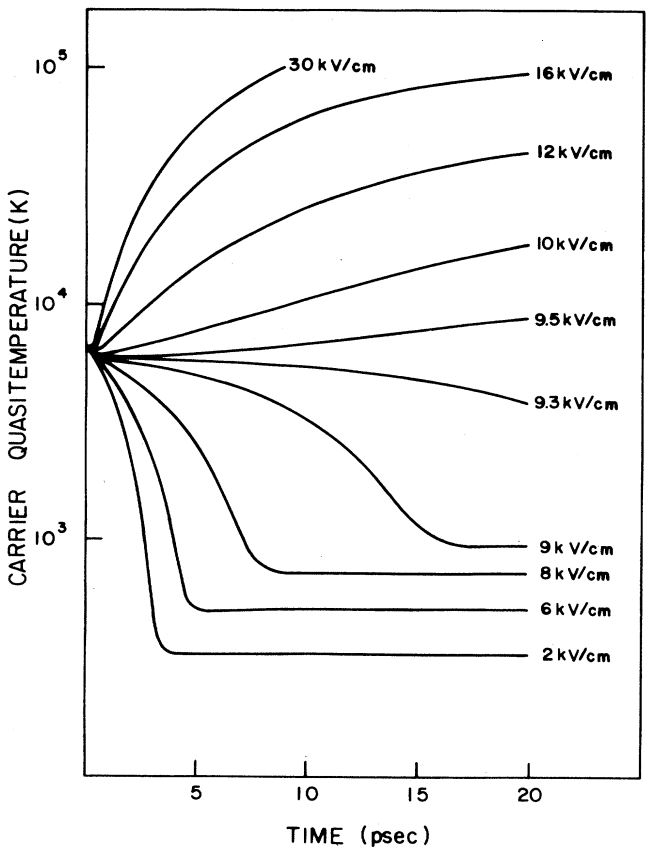

FIG. 2. Evolution of the carrier quasitemperature for several values of the electric field intensity.

shows the time evolution of the carriers' quasitemperature for several values of the electric field. It can be noted that for $\epsilon \lesssim 4$ kV/cm the carriers cool down to a steady state in less than 5 psec. For $4 \lesssim \epsilon \lesssim 9.4$ kV/cm the carriers cool down at a slower pace, and the steady state follows for delay times ranging from 5 psec to more than 30 psec. For $\epsilon \lesssim 9.4$ kV/cm the steady state follows after more than 30 psec, and the steady-state quasitemperature is larger than the initial one. This is a result of the fact that Joule heating effects overcome the energy relaxation to the phonon field. It is expected to occur starting at a value of the field intensity such that [cf. Eq. (2a)]

$$ne^2\epsilon^2\left[\frac{\tau_e}{m_e}+\frac{\tau_h}{m_h}\right]\simeq \dot{E}^{\mathrm{FR}}_{e,\mathrm{LO}}+\dot{E}^{\mathrm{FR}}_{h,\mathrm{LO}}\ . \tag{57}$$

Figure 3 shows the evolution of the LO-phonons' quasitemperature. An increase of T_{LO} of less than 10% of the reservoir temperature (300 K) can be noted, as can the fact that for fields larger than, roughly, 9.4 kV/cm, one finds smaller values of T_{LO} with increasing fields.

Figure 4 shows the evolution of the electron-drift velocity (almost identical curves are obtained for holes except that the scale of the vertical axis must be reduced by a factor of, roughly, 15).

Figure 5 displays the evolution of the momentum relaxation time and transport relaxation time; it numerically confirms the stated criterion that maxima (minima) appear when there occurs a crossover of the evolution curves of $\tau(t)$ and $\gamma^{-1}(t)$, and the latter is decreasing (increasing) at that point.

The dependence of the drift velocity in the stationary state on the electric field intensity is given in Fig. 6. This result only depends on the applied electric field intensities, being independent of the initial nonequilibrium conditions. The general form of the curve in Fig. 6, mainly the steep increase along a small interval of field intensities, has been known for a time.[16]

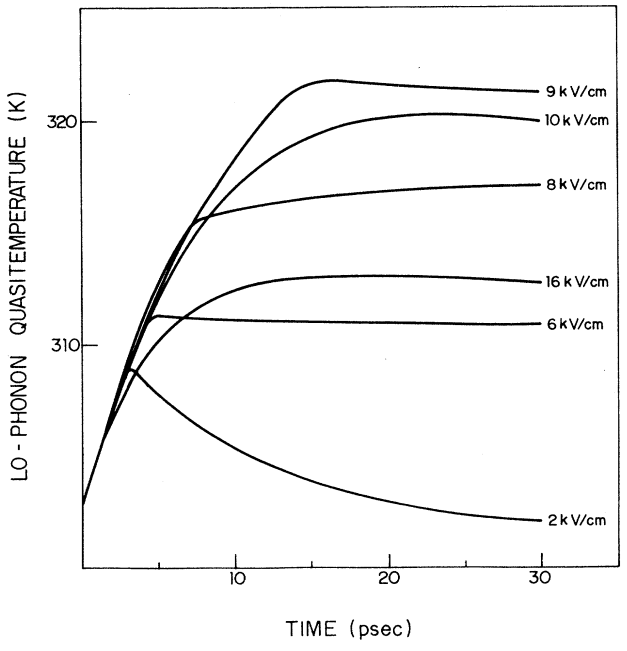

FIG. 3. Evolution of the LO-phonon quasitemperature for several values of the electric field intensity.

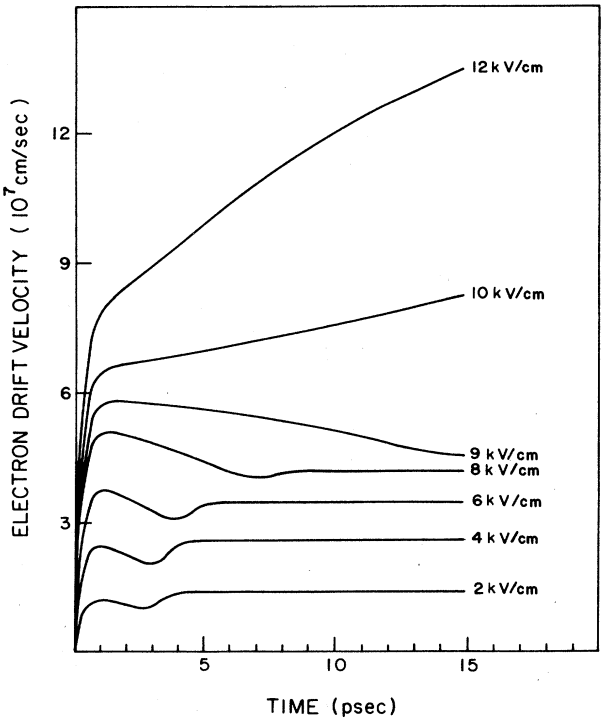

FIG. 4. Evolution of the electron drift velocity for several values of the electric field intensity.

Combination of the results shown by Figs. 5 and 6, i.e., those concerning the transient and stationary regimes, allow us to state that the ultrafast mobility of hot carriers in the central valley of GaAs (and expected to be valid for HEPS in general) has three well-defined regimes:

(i) A structured mobility with relative maximum (no overshoot) and minimum at low electric field intensities (in our case $\epsilon \lesssim 4$ kV/cm), with the mobility in the steady state following a near-Ohmic law;

(ii) a structured mobility with an absolute maximum (overshoot) and a minimum, at low to intermediate electric fields ($4 \lesssim \epsilon \lesssim 9.4$ kV/cm), and non-Ohmic behavior in the stationary state; and

(iii) normal behavior, i.e., a monotonic increase of the mobility towards its stationary value, and a near-Ohmic dependence in the latter.

We have also analyzed the effect of the initial conditions on the structured mobility transient. Figure 7 shows the evolution of the electron drift velocity for $\epsilon = 6$ kV/cm, and different values of the initial carriers' quasitemperature (i.e., increasing values of the laser frequency). At low energy transfer it follows normal behavior; with increasing energy transfer a structured mobility begins to appear, becoming more and more evident and leading, at high energy transfer, to the appearance of an increasingly more pronounced overshoot. This is the re-

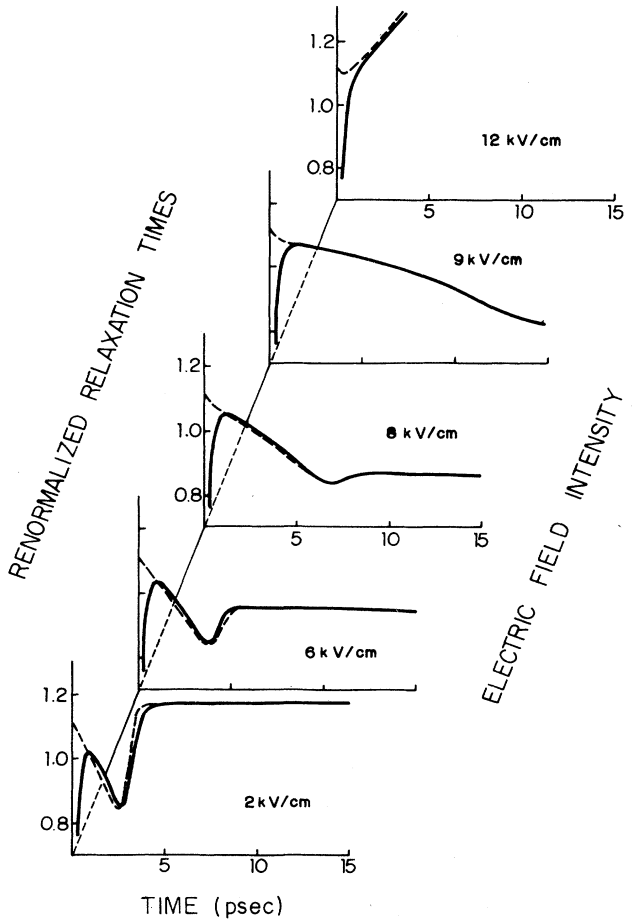

FIG. 5. Evolution of the momentum and transport relaxation times for several values of the electric field intensity.

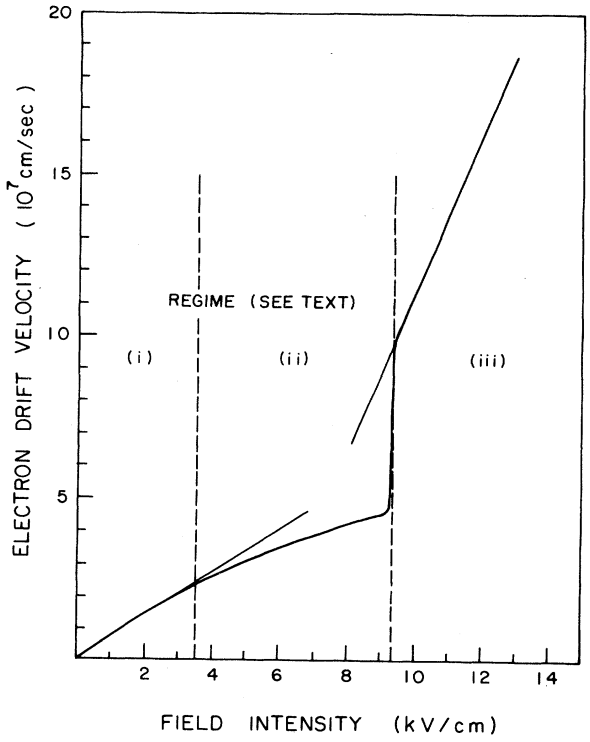

FIG. 6. Dependence of the drift velocity in the steady state on the electric field intensity.

sult of the fact that the minimum of γ^{-1} is at about $2\Theta_0$, i.e., ~ 920 K in GaAs; then, for $T_c(0)=700$ K the carrier system evolves without γ^{-1} passing through such a minimum and no structure can be produced. The subsequent values of $T_c(0)$ correspond to initial values of γ^{-1} that allow it to pass through its minimum while the carriers cool down and structure appears. With increasing values of $T_c(0)$, the initial value of γ^{-1} when on the positive slope side of the curve γ^{-1} versus T_c keeps increasing. Since the stationary value of the drift velocity is the same in all cases (independent of the initial conditions and being fixed only by the value of ϵ), for certain values of $T_c(0)$ the maximum becomes an overshoot, displaying an ever-increasing height with increasing $T_c(0)$.

Finally, we note that the existence of an instantaneous transport relaxation time allows us to write an instantaneous Einstein relation linking it to an instantaneous diffusion coefficient,

$$D_\alpha(t)=[kT_c(t)/m_\alpha]\tau_\alpha(t) . \tag{58}$$

Because of the rapid decay of the carriers' quasitemperature in the early stages of relaxation, the structure in the transient of $D(t)$ is washed out, but a very pronounced diffusion overshoot is present.

IV. DISCUSSION AND CONCLUSIONS

We have presented an analytical study of the ultrafast-mobility transient of far-from-equilibrium carriers in HEPS. For that purpose we resorted to the powerful nonlinear quantum transport theory derived from the nonequilibrium statistical operator method (NSOM) in Zubarev's approach. A coupled set of nonlinear integro-differential generalized transport equations for a basic set of nonequilibrium thermodynamic variables, deemed appropriate for the description of the macroscopic state of the HEPS, was derived. It ought to be recalled that the choice of the basic set of variables for the description of the macroscopic state of the system is one of the fundamental difficulties associated with nonequilibrium thermodynamics[17] and statistical mechanics,[3] and therefore with any approach to the NSOM.[18] In the case of HEPS one would need, in principle, to use the whole set of distribution functions for carriers and for phonons, $f_K^\alpha(t)$ and $\nu_\eta(t)$, taking as initial conditions the equilibrium ones for the latter and for the former the distributions peaked in energy space around the values compatible with energy conservation in the process of the vertical interband transition produced on absorption of one laser photon. As already noted in Sec. II, for the levels of concentration of photoinjected carriers we are considering

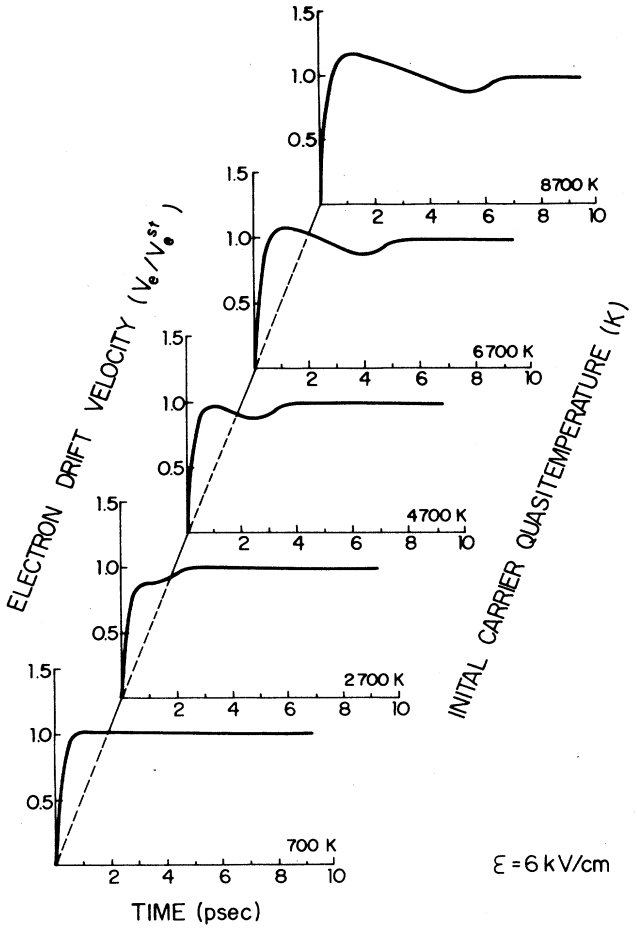

FIG. 7. Evolution of the electron drift velocity (normalized to its value in the steady state) for an electric field of 6 kV/cm, and different values of the level of photoexcitation, characterized by the initial carrier quasitemperatures.

(say 10^{18} cm^{-3} and higher), these carriers attain a condition near internal thermalization in, roughly, the 100-fsec time scale after application of the exciting laser pulse.[8] Hence, as done in the present work, after a delay time of 0.5 psec it is acceptable to work with a contracted description of the carriers' macroscopic state in terms of a quasitemperature, quasichemical potentials, and drift velocities, in an instantaneous Fermi-Dirac distribution [under high levels of excitation it goes over the expression of Eq. (15)]. Differently, the phonons do not attain a very rapid internal thermalization that allows us to describe their macroscopic state in terms of solely a quasitemperature and an instantaneous Planckian distribution, as that of Eq. (4). The phonon populations grow in time from their initial equilibrium values, while the different modes are excited by energy transfer from the carrier system, in a way that favors a certain off-center region of Brillouin zone, with internal thermalization (resulting from the interplay of electron-phonon interaction and anharmonic effects) expected to follow in the 10-psec time scale.[19] As a result we observed a slower cooling down of the carriers as compared with the one obtained using a unique quasi-temperature for all modes in each phonon branch. Consequently, the treatment we gave to the phonons in previous sections—the use of an average population characterized by $T_\eta^*(t)$—overestimates the decrease of $T_c^*(t)$ in the earlier picoseconds. As noted, the behavior of the transient mobility strongly depends on the evolution of the carriers' quasitemperature [cf. Eq. (56)], and thus the prediction concerning the structured evolution remains valid. A more precise calculation carrying the details of the evolution of the state of the phonon modes is expected to only produce numerical modifications basically—due to the slower rate of change of the carriers' quasitemperature—the appearance of the extrema at later delay times than those predicted by our theory.

These extrema, a maximum and a minimum, in the ultrafast-mobility transient of photoinjected carriers in the central valley of polar semiconductors are shown to appear at delay times during the irreversible evolution of the HEPS when there occurs a crossover of the evolution curves of the momentum relaxation time and transport relaxation time. The latter, given by Eq. (51), is a functional of the former involving memory effects, i.e., it depends on the previous history of evolution of the HEPS that determines $\gamma^{-1}(t)$. All the time dependence of the mobility (or equivalently the drift velocity) is contained in the transport relaxation time and, therefore, is governed by its evolution.

The momentum relaxation time has a minimum value for a macrostate of the HEPS between conditions of high and low excitation levels. This is so because in the first case the population of the carrier states in energy space is very low [small β_c in Eq. (5)], and the amplitude of the phonon fields does not increase appreciably (cf. Fig. 2). At low levels of excitation (low carrier quasitemperatures) conservation of energy and momentum largely reduce the number of electron states available for the scattering events. Hence, the minimum of γ^{-1} is in between: according to our calculation for a carrier's quasitemperature T_c^{ex} of nearly twice the LO-phonon's Einstein temperature. The existence of the structured mobility is totally dependent on this behavior of the momentum relaxation time. As already noted in previous sections, and corroborated by Fig. 6, for a maximum and a minimum of the drift velocity to appear, one must start with a sufficiently photoexcited HEPS (initial T_c larger than T_c^{ex}), followed by a cooling of the carrier system until values of T_c are smaller than T_c^{ex}.

With increasing field intensities, T_c decreases with a slower pace and its value at the steady state, T_c^{ss}, is larger and larger, accompanied by a mobility smaller and smaller, with the result that the value of γ_{ss}^{-1} approximates its minimum value. On the other hand, the height of the maximum of τ is weakly altered since it depends exclusively on the evolution of γ^{-1} towards its minimum value. With increasing values of the electric field intensity and decreasing values of γ_{ss}^{-1} ($=\tau_{ss}$), Joule heating effects become strong and overcome the process of energy

relaxation to the lattice [cf. Eq. (57)]. Consequently, T_c^{ss} in the neighborhood of T_c^{ex} ($\sim 2\Theta_0$ or, roughly, 1000 K in GaAs attained at ~ 9 kV/cm) is strongly dependent on the electric field intensity, and so is γ_{ss}^{-1}, leading to a non-Ohmic dependence of the drift velocity. Then, the regime with overshoot is strongly correlated with the region of values of electric field intensity for which the stationary mobility is non-Ohmic as confirmed by inspection of curves 4 and 6.

Summarizing, for a sufficiently intense electric field the carrier system keeps heating up (or starts to cool down and next heats up) so that the momentum relaxation time does not attain its minimum and structured mobility is excluded (In Fig. 5 the case for 12 kV/cm). Hence, there exist a maximum value of the field above which the mobility only presents normal evolution. This regime follows, on increasing field strength, from another one where structured mobility is present and the maximum is an overshoot (In Fig. 5, the cases for 6, 8, and 9 kV/cm). The height of the overshoot diminishes with decreasing field intensity and there is a lower limit of this intensity below which this maximum is no longer an overshoot (In Fig. 5 the case for 2 kV/cm). Further, as shown by Fig. 7 and the ensuing discussion in Sec. II, there is a lower level of photon laser energy for the phenomenon to occur, viz., the one that allows excess kinetic energy of the carriers to be high enough for the macroscopic state of the system to allow, at the start, for the momentum relaxation time to decrease with decreasing temperature.

Experimental observations of ultrafast mobility in HEPS are scarce, and the existing few are not detailed enough.[20,21] It may be mentioned that there is a certain qualitative and semiquantitative agreement with Hammond's measurements[21] in that he reports an inferior and superior limit of the field intensity for an overshoot to be observed. Also, it must be stressed that we have studied the dependence of the mobility of nonequilibrium carriers in HEPS in a single valley. However, the band structure of direct-gap polar semiconductors displays multiple valleys, and therefore intervalley scattering of carriers needs to be considered. It could lead to additional structure at sufficiently high levels of excitation as a result of the transference of carriers to higher energy valleys where they have larger effective masses, but this effect seems to be smoothed out by carrier collisions.[22]

Concerning Zubarev's method used in this work, as noted we applied it to the study of relaxation effects on optical properties of HEPS;[10,23,24] its seemingly first application to the study of transport properties in solid-state systems is due to Kalashnikov,[24] more recently it was discussed in this connection by Ferry *et al.*,[25] and Liu *et al.*[26] applied it to the study of the steady state of high-field electron transport in multivalley semiconductors.

ACKNOWLEDGMENTS

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