A theoretical analysis of the picosecond and sub-picosecond infrared-absorption spectroscopy of hot holes in germanium

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We present an ensemble Monte Carlo study of the thermalization and relaxation of photoexcited heavy and light holes in gallium doped p-type germanium for doping densities of $1 \times 10^{16}$ and $7 \times 10^{16}$ cm$^{-3}$. We analyze recently published data obtained by picosecond excite-and-probe absorption spectroscopy in the mid infrared, i.e., with excitation energies of 124 and 136 meV, and at lattice temperatures of 30, 80, and 300 K. The Monte Carlo simulations contain all types of intraband and interband hole–hole and hole–optical–phonon scatterings, as well as free-carrier-induced ionizations of neutral acceptors and captures of free holes by ionized acceptors. Good agreement is found between the simulated and the measured absorption changes as functions of time. Quite generally it turns out that the initially created nonthermal features in the heavy- and light-hole distribution functions persist to times of typically twice the pump–pulse duration, thereby noticeably contributing to the initial nonlinear response. For longer times, the practically thermalized distributions give absorption changes which exclusively reflect the cooling dynamics of the heavy holes. © 1997 American Institute of Physics.

I. INTRODUCTION

The time-resolved spectroscopy of photoexcited electrons e and holes h has, for many years, provided the most detailed information on the dominant scattering mechanisms of "hot" carriers in semiconductors. The standard experimental scenarios in the visible or near infrared involve e–h excitations across the principal bandgap, and due to the strong intercarrier coupling it is difficult to directly extract information on the relaxation dynamics of the separate electron and hole subsystems. In this context, the infrared spectroscopy of photoexcited holes in a p-doped material has the advantage of providing information on the intravalence and intervalence band processes of holes without the participation of electrons.

The experimental data underlying the present theoretical study were recently published by Wörner, Elsässer, and Kaiser (henceforth cited as WEK); they were obtained by mid-infrared picosecond pump-and-probe absorption spectroscopy of gallium doped p-type germanium for sample thicknesses of 0.2 and 0.04 cm, doping densities of $1 \times 10^{16}$ and $7 \times 10^{16}$ cm$^{-3}$, and lattice temperatures of 30, 80, and 300 K. Figure 1 gives a schematic picture of the experimental WEK scenario. The photon energy $h\nu_{\text{pump}}$ of the pump pulses was 124 or 136 meV, with pulse energies of 5 $\mu$J. The duration and bandwidth of the pump-and-probe pulses were 2 ps and 1.8 meV, respectively. Depending on the sample and the lattice temperature, the corresponding (calculated) excitation densities $n_{\text{exc}}$ ranged from $4 \times 10^{15}$ to $3 \times 10^{16}$ cm$^{-3}$. With the above choice of photon energies, the experiments allowed us to study photo- and scattering-induced transitions of holes between the heavy-hole (hh) and the light-hole (lh) valence band, as well as free-carrier-induced acceptor ionizations and the corresponding hole-capture processes, without contributions of the split-off valence band. Except for one example, our following theoretical analysis will be restricted to the case of degenerate pump and probe frequencies, i.e., to $h\nu_{\text{probe}}=h\nu_{\text{pump}}$.

In Sec. II we present the essentials of our ensemble Monte Carlo (EMC) simulation of the coupled hole-acceptor-phonon system and discuss the underlying transport description of the photoexcitation and the scattering processes. In Sec. III the results of the simulations are compared with the experimental data. The concluding Sec. IV relates the results to a long standing controversy about the possibility of strongly nonthermal hole distribution functions in this type of experiments.

II. THEORY

As the EMC technique and its application to semiconductor transport have been extensively described in the literature, we only emphasize the technicalities specific to the present analysis. The simulation starts with the equilibrium carrier ensemble of the doping background just before the onset of the excitation pulse. Thereafter the photogenerated holes and corresponding hole vacancies are added according to the temporal pulse shape (assumed to be gaussian) and the experimentally determined excitation spectrum. For the experimental peak intensity $I_p=30$ MW/cm$^{-2}$ of the pump pulse and by assuming relative reflection losses of 2/3 and a quantum efficiency per absorbed photon of 1, the excitation spectrum was obtained by use of an effective pulse spectrum proportional to the measured linear absorption profile. The hh–lh profile was well approximated by a superposition of two gaussians.
The second contribution to Eq. 1 describes transient changes of the hh and lh state occupations. The first line of Eq. (1) perturbation theory, 6,7 distribution functions f_{hh}(k,t) and f_{lh}(k,t) through standard perturbation theory, 6,7

$$I(\nu) = \frac{1}{1+0.33/\cos h(12w_2)} \frac{1}{\cos h[w_1(h\nu-h\nu_0)]}$$

with the nominal excitation energy $h\nu_0$ (124 or 136 meV), respective widths $h\Delta \nu_1 = 10$ meV and $h\Delta \nu_2 = 22$ meV, and $w_2 = 2.6339/(\Delta \nu_2)$. In this way, effects of band warping could be fully included in the simulated photoexcitation spectrum.

The spectrally and temporally resolved absorption coefficient was obtained from the simulated hh and lh k-space distribution functions $f_{hh}(k,t)$ and $f_{lh}(k,t)$ through standard perturbation theory. 6,7

$$\alpha(\nu,t) = \frac{8\sqrt{2}e^2}{m_0^2\hbar^2nc} \frac{m_{hh}m_{lh}}{m_{hh}-m_{lh}}$$

$$\times A_{12} \left[ \left( \frac{m_{hh}m_{lh}}{m_{hh}-m_{lh}} \right)^{3/2} \sqrt{h\nu} \left[ f_{hh}(k,t)-f_{lh}(k,t) \right] \right]$$

$$+ \frac{16\pi m_{lh}^{3/2}}{(1+2a_0^2m_{lh}E_{hh}/h^2)^4}$$

$$\times n_{\text{dop}} \left[ f_{\text{dop}}-f_{hh}(k,t) \right] \sqrt{h\nu-E_{\text{ion}}},$$

where $e$ denotes the elementary charge, $m_0$ the free electron mass, $c$ the light velocity, $n$ (4) the refractive index, $E_{\text{ion}}$ (= 10 meV) the ground-state ionization energy of the gallium acceptor, $a_0$ the corresponding effective Bohr radius, $E_{lh}$ the lh excitation energy resulting from an acceptor ionization, and $k$ the number of the optically coupled states. The first line of Eq. (2) describes the hh–lh interband excitation, where the difference $f_{hh}(k,t)-f_{lh}(k,t)$ contains the effects of the pump-pulse-induced transient changes of the hh and lh state occupations. The second contribution to Eq. (2) describes the transitions from neutral acceptors into the hh band, where $f_{\text{dop}}$ denotes the relative population $n_{\text{A}}^{(0)}/n_{\text{dop}}$ of neutral acceptors within the total doping background. Various values for the parameter $A_{12}$, arising from the directional average of the optical transition amplitude, can be found in the literature; we have used its most recently proposed value of 7.2. 8

The scattering dynamics contained (i) all types of intraband hole–hole and hole–optical–phonon scattering, (ii) inter-valence-band scattering induced by both hole–optical–phonon and hole–hole scatterings, 9,10 and (iii) free-carrier-assisted ionizations of neutral acceptors and captures of free holes by ionized acceptors. 5,10 Expecting only a secondary influence of band warping or band nonparabolicities on the overall scattering dynamics, the simulation of the scattering processes was performed for quasifree heavy and light holes, with average effective masses $m_{hh}/m_0 = 0.36$ and $m_{lh}/m_0 = 0.04$. We checked the assumption of a parabolic lh band by keeping track of the relative number of simulated light holes reaching the regime of strong nonparabolicity (at energies beyond 170 meV), and found a negligible fraction. Similarly, and again for reasons of computational simplicity, only the acceptor (1s) ground state, with an ionization energy $E_{\text{ion}} = 10$ meV, was considered in the total and differential cross sections for capture and ionization (see Appendix A).

In view of the extreme numerical expenditure needed for the implementation of a fully dynamical free-carrier screening of the interparticle Coulomb interaction, we used a recently developed and computationally practicable nonequilibrium version of the long-wavelength limit of the static random phase approximation (RPA) susceptibility (see Appendix B). Pauli blockings and overlap factors were included by standard Monte Carlo rejection techniques. 3–5

For both the intraband and interband carrier-optical-phonon scatterings we used a value of 6.3 eV/Å for the optical deformation-potential constant. 2 Acoustic-phonon scatterings and elastic interband scatterings at ionized acceptors, although initially included in our scattering scenario, turned out to be negligible and were omitted in all later simulations.

The simulated carrier ensembles contained typically 20 000 particles. For the comparison between theory and the measured absorption changes the simulated nonlinear optical response was finally convoluted with the spectral and temporal shape of the probe pulse; here we assumed identically shaped pump-and-probe pulses.

III. RESULTS

The experimental data had been taken two samples; sample A with doping density $n_{\text{dop}} = 1.1 \times 10^{16}$ cm$^{-3}$ and a thickness $d$ of 0.2 cm, and sample B with $n_{\text{dop}} = 7 \times 10^{16}$ cm$^{-3}$ and $d = 0.04$ cm.

The basic experimental quantity for the degenerate pump-probe absorption spectroscopy of our concern is the negative absorption change (or “bleaching”), defined in terms of the pump-pulse-induced relative transmission change $\Delta T(t) = T(t)/T_0$ through $1^2$ $\Delta A(t) = -\ln[\Delta T(t)]$. Neglecting the frequency dependence of the reflection losses, one thereby obtains

$$\Delta A(t) = -[\alpha(t)-\alpha_0]d$$

as the basic theoretical quantity with which to compare the experimental data.
Figure 2 shows the experimental absorption change $\Delta A(t)$ (squares) and the results of the EMC simulation (full line) for sample A and a lattice temperature of 300 K over a time interval of 15 ps; time zero taken at the maximum of the excitation pulse. The corresponding Fig. 3 shows the calculated hh distribution function $f_{hh}$ and the lh distribution function $f_{lh}$ versus energy at various time instants on a semilogarithmic plot. The full vertical lines mark the energetic positions of the optically coupled hh and lh band states, and the dashed vertical lines the position of the additional acceptor-lh excitation channel. An internally thermalized distribution, i.e., a Maxwellian for the present range of temperatures and carrier densities, would yield a straight line. However, we note a pronounced peak (state filling) in $f_{hh}$ and a corresponding dip (state depletion) in $f_{lh}$ at the hh–lh excitation energies throughout the excitation pulse. This means a reduced absorption rate ($\propto f_{hh}^{-} f_{lh}$) and a resulting bleaching. In view of the fact that the calculations contained no adjustable parameter, the agreement between theory and experiment in Fig. 2 is very satisfactory. To provide an estimate of the additional role of carrier heating and cooling, the dashed curve in Fig. 2 shows the result obtained by using exact Maxwellians with a carrier temperature $T_c$ determined by the simulated mean particle energy of the total, i.e., hh plus lh, carrier system via $\langle E \rangle_{hh+lh} = 3k_BT_c/2$. Within this approximation, the qualitative shape of $\Delta A(t)$, with its maximum around $t=0$, could be understood as the initial increase of $f_{hh}^{\text{max}}(k,t)$ during the laser-induced heating of the total carrier system and its ensuing decrease due to energy losses, mainly via optical–phonon emissions. However, the resulting bleaching is much too small, providing strong evidence for the decisive role of the nonthermal features in the hh and lh distributions. As a further technical point we should note that, in any case, the use of a common carrier temperature for heavy and light holes could not be justified in view of the pronounced difference between the simulated mean hh and lh energies shown in Fig. 4. Insisting on the use of a single carrier temperature, an attempt to use $T_c$ as parameter to fit the experimental data would require unrealistically high hole temperatures, e.g., $T_c \sim 1000$ K at $t=0$, which would correspond to a mean particle energy of 130 meV as compared to the simulated mean particle energies in Fig. 4. We therefore conclude that the experimental results of Fig. 2 can only be explained by the combined effects of nonthermal hh and lh distributions and of carrier heating and cooling.

FIG. 2. Absorption change $\Delta A$ as a function of time during and after the pump pulse for sample A; the experimental parameters are listed in the figure.

FIG. 3. Distribution functions of heavy and light holes at various time instants for the experimental parameters of Fig. 2.

FIG. 4. Mean kinetic energy of light and heavy holes as a functions of time for the experimental conditions of Fig. 2.

FIG. 5. Negative absorption change $\Delta A$ as a function of time during and after the pump pulse for sample A; the experimental parameters are listed in the figure.
A qualitatively different time evolution of $\Delta A(t)$ is found for the same excitation conditions, but for a lattice temperature of 80 K. Figure 5 again shows a bleaching at early times, but a transition to induced absorption at times beyond 5 ps. Also for this case, the corresponding carrier distributions, shown in Fig. 6, directly reflect this experimental feature. For times between 2 and 5 ps, state filling of light holes and state depletion of heavy holes cause a transient initial bleaching. At later times, however, after the end of the pump–pulse, the distributions are already practically thermalized, with a common carrier temperature $T_c$ higher than the lattice temperature $T_L$. This causes $f_{hh} < f_{lh}$, and consequently the absorption to be larger than at equilibrium and leads to a negative $\Delta A(t)$, again in satisfactory agreement with experiment (Fig. 5). The eventual decrease of the induced absorption is caused by the cooling of the heavy holes by phonon emissions, in agreement with the original WEK interpretation. The dashed line in Fig. 5 again shows the result obtained by use of a common carrier temperature corresponding to the simulated mean particle energy. Here, the $T_c$ model even fails to yield the short-time bleaching. Naturally, the agreement of the results of the $T_c$ model and the detailed EMC result at times beyond the thermalization regime is much better. The remaining disagreement between the theoretical results can be explained by the fact that the simulated distributions are not yet completely thermalized. Indeed, the calculated mean hh and lh energies still differ by a few percent at times $t > 15$ ps.

The analogous results for the higher-doped sample B are qualitatively and quantitatively similar to our foregoing examples. The reason is that $\alpha(t) \propto f_{hh} - f_{lh} \propto n_{\text{dop}}$ is approximately six times larger, but $\Delta A(t)$ is practically unchanged due to the five times smaller sample thickness $d$. Figures 7 and 9 again show a semiquantitative agreement between experiment and theory. (The reader is also referred to Figs. 8 and 10.) The remaining discrepancies, in particular around the bleaching maxima, may arise partly from our neglect of the nonsphericity (warping) of the bands and partly from our neglect of excited acceptor states.

Figures 11 and 12 finally present results for sample B at $T_L = 30$ K for the case of $h\nu_{\text{pulse}} = 136$ meV, $h\nu_{\text{pulse}} = 130$ meV, and $I_L = 20$ MW/cm$^2$. Both experiment and theory yield induced absorption at all times, which can again be explained by the time evolution of the distribution functions, as shown in Fig. 12. At this low temperature, the equilibrium value of $f_{hh}$ at the hh–lh excitation energy is practically zero...
due to carrier freezeout. So the lh excitation proceeds almost exclusively through the acceptor–lh channel, as seen from the dominant lower-energy excitation peak (around the dashed vertical line) of $f_{lh}$ in Fig. 12. The initially small population of the heavy-hole band is soon increased and heated by lh → hh band transfers and lh–hh collisions, resulting in an increase of $f_{hh}$ in the hh–lh absorption regime. It turns out that this effect overcompensates the bleaching action of the lh state filling in the acceptor–lh channel. As the dominant contribution to $\Delta \sigma(t)$ is now the lh–induced heating and the ensuing phonon-induced cooling of the heavy holes, the result obtained within the $T_c$ description by use of the simulated mean particle energy (dashed line) is now at least in qualitative agreement with the EMC result (full line). It is worth noticing that possible effects of impurity banding cannot be excluded at these low temperatures; but any corresponding modification of the lh-excitation dynamics would not call into question the dominant role of the hh-heating and cooling mechanism for the low-temperature data of Fig. 11.

IV. CONCLUSIONS

Two features were shown to determine the transient absorption changes observed in recent mid-infrared picosecond pump-probe experiments on p-doped germanium at moderate doping densities: (i) state-filling within the light-hole band and state-depletion in the heavy-hole band and (ii) energy transfers to and from the hh subsystem. The relative importance of these two effects strongly depends on the lattice temperature and the detailed excitation and detection conditions. In all simulations it turned out that, independent of lattice temperature and doping density, the hole distributions are non-Maxwellian for times of several picoseconds around the maximum of the 2 ps excitation pulse. Only at later times are the hh and lh distributions well described by Maxwellians.

We believe that these findings are of relevance to a long-standing controversy about the possibility of strongly nonthermal hole distributions in ir excited p-Ge, a question of some relevance for a more detailed understanding of the "hot-hole" laser. Assuming that the carrier–carrier scatterings were fast enough to establish an internal and mutual hh and lh thermalization at times smaller than the experimental time resolution, WEK originally interpreted their data as reflecting the cooling dynamics of a completely thermalized hh–lh system. Indeed, the data could be well reproduced by the standard energy-loss rates within the conventional carrier-temperature model by use of an appropriately fitted time-dependent hole temperature. This agreement between the data and the carrier-temperature model was taken as strong indication against the possibility of nonthermal hole distributions in infrared-excited p-Ge. Such strongly nonthermal features, like spectral-hole burning, had earlier been postulated to explain the results of far infrared studies of photoexcited hh–lh systems in p-Ge on the nanosecond time scale, i.e., under quasisteady-state conditions. In this context our results demonstrate that both effects, i.e., state-filling and depletion as well as heating and/or cooling of carrier distributions, can decisively contribute to the transient optical response, at least under the conditions of the WEK experiments.
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APPENDIX A: RATES FOR HOLE-ACCEPTOR TRANSITIONS

In the following we shall summarize the essential steps in the derivation of the differential and total rates for hole–hole scattering induced ionization and capture processes, indicating also some computational details for their implementation in the EMC simulations. The interaction of a free hole with either a 1 s electron in a neutral acceptor or an ionized acceptor involves transition amplitudes of the basic type

\[ M_{12} = \langle 1s, k_1 | H_{\text{int}} | k_2 \rangle. \]  

(A1)

With plane waves for the free-carrier states and standard 1 s wave functions and with quasistatically screened Coulomb interactions, the spatial integrations in Eq. (A1) yield

\[ M_{12} = \frac{4 \pi e^2}{\epsilon_0 (k_2 + q_s)^2} \left( \Omega a_0 \right)^{3/2} \frac{1}{a_0} \left[ \frac{1}{(1/a_0)^2 + (k_2)^2} \right] \times \frac{1}{\left[ ((1/a_0)^2 + (k_2 + q_s)^2) \right]}, \]  

(A2)

where \( \Omega \) denotes the crystal volume, \( q_s \), the inverse screening length, and \( k = k_1 - k_2 \). To obtain the total scattering rate, Fermi’s golden rule and lengthy, but straightforward integrations over \( k_2 \), and the angular part of \( k_1 \), finally lead to

\[ I_{\text{tot}}^{\text{ion}}(k_1) = \frac{2 e^4 m_2 n_{A} (1)}{|h^2 e_0 a_0|} \int k_1 d k_1 (k_1) f_j(Q(k_1)) \]

\[ \times \left[ \frac{2 Q I_1}{(p^2 + Q^2)^s} + \frac{I_2}{(p^2 + Q^2)^s} \right], \]

(A3)

and

\[ I_{\text{tot}}^{\text{cap}}(k_1) = \frac{2 e^4 m_2 n_{A} (1)}{|h^2 e_0 a_0|} \int k_1 d k_1 (k_1) f_j(Q(k_1)) \]

\[ \times \left[ \frac{2 Q I_1}{(p^2 + Q^2)^s} + \frac{I_2}{(p^2 + Q^2)^s} \right], \]

(A4)

with the following abbreviations:

- ionization:
  \[ k_i = 0, \quad k_u = \frac{1}{h} \sqrt{2 m_i (E_1 - E_{\text{ion}})}, \]
  \[ Q^2 = \frac{m_i}{m_i} \left( k_i^2 - k_{i'}^2 - \frac{2 m_i E_{\text{ion}}}{h^2} \right), \]

- capture:
  \[ k_i = \frac{1}{h} \sqrt{2 m_i (E_1 + E_{\text{ion}})}, \quad k_u = \infty, \]
  \[ Q^2 = \frac{m_i}{m_i} \left( k_i^2 - k_{i'}^2 - \frac{2 m_i E_{\text{ion}}}{h^2} \right), \]

where \( n_{A} \) and \( n_{A}^{-} \) denote the density of neutral and ionized acceptors; the indices \( i \) and \( j \) refer to the partner particle \((k_i \rightarrow k_{i'})\) and to the captured or released electron, respectively. For both types of scatterings we have defined

\[ I_1 = \frac{2}{(k_i^2 + k_{i'}^2 + q_s^2)^2 - 4 k_i^2 k_{i'}^2}, \]

\[ I_2 = \frac{\alpha}{6 k_i k_{i'}}, \quad J_1 = \frac{\beta}{6 k_i k_{i'}}, \quad J_2 = \frac{\gamma}{6 k_i k_{i'}}, \quad \left[ \frac{3 b}{c} J_{22} + \frac{5}{c} J_{21} + \frac{1}{c} \left( \frac{1}{x_+ X_+} - \frac{1}{x_- X_-} \right) \right], \]

\[ I_3 = \frac{2 Q}{k_i k_{i'}}, \quad b = 2 (Q^2 + q_s^2 - b^2), \]

\[ c = (p^2 + Q^2)^2 + q_s^2 (2 Q^2 - p^2), \]

\[ \alpha = 6 Q, \quad \beta = 12 Q q_s^2 - 12 Q p^2 - 20 Q^3, \]

\[ \gamma = 6 Q p^4 + 6 Q q_s^2 + 6 Q^5 - 20 Q^3 q_s^2 + 12 p^2 Q^3 - 12 p^2 q_s^2, \]

\[ \Delta = 4 c - b^2 = 16 p^2 Q^2, \quad X_{\pm} = x_+^2 + b x_{\pm} + c, \]

\[ T_{\pm} = \frac{1}{\sqrt{\Delta}} \arctan \frac{2 x_{\pm} + b}{\sqrt{\Delta}}. \]

The final integration over \( k_{i'} \) in Eqs. (A3) and (A4) has to be performed numerically. The total scattering rates, Eqs. (A3) and (A4), turn out to be only weakly energy dependent; at room temperature and for a doping density of...
They are of the order of 0.1 event per ps. Following the conventional Monte Carlo strategy, the final wave number of particle 1 after the scattering is determined by randomly selecting one of the equally distributed \( k_1 \), values out of the respective integration range and using the corresponding probability distribution in Eqs. (A3) and (A4). The latter are given by the \( k_1 \)-dependent part of the differential scattering rate, i.e.,

\[
P(k_1 \rightarrow k_1') = k_1^2 \left( \frac{2Q_1}{(p^2 + Q^2)^{3/2}} + \frac{\tilde{I}_2}{(p^2 + Q^2)^{1/2}} \right)
\]

(A5)

for an ionization and

\[
P(k_1 \rightarrow k_1') = k_1^2 f_3 Q(k_1') \left[ \frac{2Q_1}{(p^2 + Q^2)^{3/2}} + \frac{\tilde{I}_2}{(p^2 + Q^2)^{1/2}} \right]
\]

\[
+ \frac{\tilde{I}_3}{(p^2 + Q^2)^{1/2}}
\]

(A6)

for a capture process.

Given \( k_1 \) and the final wave vector \( k_1' \) of the simulated particle, the energy of the partner particle (of type \( j \)) is fixed. In case of a capture one removes the energetically nearest such particle from the ensemble and in case of an ionization one adds a particle with arbitrary direction of its wave vector, in both cases with the appropriate changes of the densities \( n_A^{(0)} \) of neutral and \( n_A^{(-1)} \) of ionized acceptors.

**APPENDIX B: FREE-CARRIER SCREENING**

The screening of the carrier–carrier and the long-range polar carrier–phonon interactions by free carriers should in principle be taken as dynamical and self-consistent, i.e., the screening model should account for the inelasticity of the scatterings and for the nonequilibrium and multicomponent nature of the free carrier plasma. The following version of the long-wavelength limit of the static RPA susceptibility has been recently found to be in good agreement with the results of molecular-dynamical simulations of photogenerated hot-carrier plasmas. In this screening model the nonequilibrium distributions of the heavy and light holes are consistently taken into account by use of instantaneous ensemble averages over the inverse single-particle energies. In an obvious notation, the expression for the square of the corresponding inverse screening length then reads

\[
q_s^2 = \frac{4 \pi e^2}{\varepsilon_\infty} \frac{1}{2} \sum_{i=\text{hh},\text{lh}} n_i \left( \frac{1}{E_i} \right).
\]

(B1)

For c–c scattering an additional quasidynamical retardation effect was introduced by allowing only particles of mass \( \approx m_\alpha \) to screen a carrier–carrier interaction between two particles of masses \( m_\alpha \) and \( m_\beta \approx m_\gamma \).

The remaining long-range degrees of freedom within the hh–plus–lh plasma are treated in analogy to the standard procedure for the equilibrium electron gas by introducing two-component-plasma modes and the corresponding carrier-plasmon interactions. However, the plasmon scatterings turned out to negligibly contribute to the net energy-transfer rates and were therefore omitted in most simulations.