Theoretical study of sub-to-picosecond laser pulse interaction with dielectrics, semiconductors and semiconductor heterostructures

T Apostolova

Institute of Nuclear Research and Nuclear Energy,
Bulgarian Academy of Sciences, 72 Tzarigradsko Chaussee, 1784 Sofia, Bulgaria

E-mail: tzveta@mail.inrne.bas.bg

Abstract. A summary of theoretical investigations of short pulse (τ ≤ 1 ps) laser interaction with dielectrics, semiconductors and semiconductor heterostructures is presented. The time-dependent kinetic Fokker-Planck type equation for excited conduction electrons is used to describe this interaction in SiO₂ and is systematically derived for GaAs. The energy spectra of the electron distribution function and the time dependence of the electron density are presented to illustrate the role of the different physical processes responsible for the conduction electron dynamics. Some possible types of laser induced damage in dielectrics and semiconductors including optical, electrical and structural damage are explored. In addition, the photon-absorption-induced intrasubband and intersubband phonon scattering of conduction electrons in GaAs/AlGaAs quantum wells are predicted by using the Boltzmann equation in the presence of a normally incident mid-IR pulsed laser field.

1. Introduction

Ultra-short (picosecond to femtosecond) laser pulses deliver very high peak power and can be produced by small, light, portable devices and by free electron lasers (FEL). The applications of these increasingly available ultra-fast lasers are limited by the laser damage to optical components. Moreover, the sensitive semiconductor and semiconductor heterostructure-based photodetector devices degrade in performance when irradiated with intense ultra-short pulse laser field.

Compared to continuous lasers, ultra-short pulse laser systems damage targets by localized effects before the heated electrons produced during the interaction have time to diffuse away. For that reason damage caused by ultra-short laser pulses is characterized by a minimum of collateral damage and is controllable on a microscopic scale so there is a large number of potential applications such as micromachining of optical materials (e.g. producing arrays of damage dots for optical memories), microlithography, medical surgery, etc. for which the controlled deposition of laser energy in solids is crucial. These are the positive aspect of the laser induced breakdown (LIB) of materials.

The drawbacks and the distinctive advantages of the ultra-short pulse laser induced breakdown mentioned above require fundamental understanding of the laser-induced damage mechanisms in specific materials in particular and in solids in general.

The goal of this paper is to present a summary of theoretical investigations of picosecond-to femtosecond laser pulse interaction with selected dielectric and semiconductor materials describing the occurring microscopic processes. The choice of the investigated materials is governed by their
practical applications in optical and semiconductor technologies which consequently determines the used laser wavelengths and pulse durations.

Some results, not previously published in peer reviewed journals are presented which can be useful in creating technologies for laser processing of materials and fabricating components of applied quantum electronics.

The governing microscopic processes occur within the time scale of the laser pulse so that the appropriate dynamical models have to explicitly account for them. The kinetic Fokker-Planck equation is one way to describe the physics of the processes taking place when a dielectric or a semiconductor gets irradiated by intense laser field and can be used to explain how they lead to a laser-induced breakdown of these materials. The damage due to high intensity, ultra-short pulse laser fields is due to the production of huge amounts of hot electrons which form electron plasma leading to the breaking of chemical bonds, and the creation of defects inside the target materials.

2. Theoretical picture
In the theoretical picture, LIB of dielectrics and semiconductors directly results from electron production in the conduction band (CB) by photoionization via coherent absorption of photons from the incident laser field followed by incoherent interband Coulomb scattering through induced collisional ionization with excited conduction electrons. Collisional ionization can occur if electrons in the CB gain energy until they reach a threshold kinetic energy equal to the band-gap. These excited electrons then may collide with valence electrons and promote them into the CB by energy exchange. The electrons excited by collisional ionization in the CB can in turn reach threshold energy and excite more valence electrons to the CB which leads to avalanche ionization. In order to reach the threshold energy for collisional ionization, conduction electrons have to gain energy both classically and quantum-mechanically from the incident laser.

For the study of the laser-induced damage in dielectrics we have adopted the theoretical approach used by Hollway and Fradin [2] and later extended by Stuart et al [3] and further have improved it by including additional terms in the Fokker-Planck type kinetic equation. To investigate the laser-induced damage in semiconductor materials we derive a generalized Fokker-Planck quantum dynamical theory that is a first-order approximation of a general semi-classical Boltzmann theory. We also calculate the hot electron distribution function in the first and second subband of quantum wells irradiated by a normally incident mid-infrared pulsed laser field by including the photon-assisted phonon-scattering process in a Boltzmann equation for phonon energies that are smaller than the energy separation between two electron subbands in the quantum well. This electron distribution can be used to evaluate the nonequilibrium electron temperature and the average electron energy which are essential in understanding the laser damage of semiconductor materials.

By analyzing our numerical results we gain insight into the microscopic mechanisms for laser-induced damage in dielectric optical components and semiconductor devices.

In the following sections we use the notation pertaining to the relevant cited papers.

3. Laser interaction with dielectrics
A Fokker Planck equation with a complete set of rates is used to describe the free electron production in the wide band gap dielectric fused silica in terms of the energy distribution function \( f(\epsilon,t) \) when the material is subjected to intense laser field. The conditions under which the kinetic equation is derived and its range of validity are detailed in [1], [4]. The kinetic equation for \( f(\epsilon,t) \) is given by

\[
\frac{\partial f(\epsilon,t)}{\partial t} + \frac{\partial}{\partial \epsilon} \left( R_i(\epsilon,t)f(\epsilon,t) - U_i(\epsilon)f(\epsilon,t) - D(\epsilon,t) \frac{\partial f(\epsilon,t)}{\partial \epsilon} \right) = S(\epsilon,t),
\]

where \( f(\epsilon,t) \) is the number density of electrons with kinetic energy between \( \epsilon \) and \( \epsilon + d\epsilon \) at time \( t \). All the coefficients in equation (1) are detailed in [1]. The joule heating given by \( R_i(\epsilon,t) = \frac{1}{3} \sigma(\epsilon)E^2(t) \) and the diffusion in energy space \( D(\epsilon,t) = \frac{2}{3} \sigma(\epsilon)E^3(t) = 2xR_i(\epsilon,t) \) are functions of the laser intensity, and
account for the electron energy gain from the laser field in terms of the electron conductivity
\[ \sigma_e(t) = \left[ e^2 \gamma_e(t) \right] m^* \left[ 1 + \omega^2 \tau_e^2 \right] \]
where \( m^* \) is the effective electron mass, \( e \) is the electron charge and \( \gamma_e(t) = \tau_e \left[ \tau_e(t) \right] \) is the momentum scattering rate. The electron energy loss to phonons is
\[ U_i(e) = h \omega_{\rho k} \gamma_i(e), \]
where \( \gamma_i(e) \) is the rate of electron energy transfer to the lattice and \( h \omega_{\rho k} \) is the characteristic phonon energy. The net number of electrons per unit volume whose energy increases from a value less than \( \epsilon \) to a value greater than \( \epsilon \) per unit time is the electron charge and is defined as the current in energy space. The term \( S(e, t) = R_{mp}(e, t) + R_{rec}(e, t) + R_{\rho}(e, t) \) contains all the sources and sinks of electrons given in detail in [1]. The impact ionization term
\[ R_{mp}(e, t) = -v_i(e)f(e, t) + 4\nu_e(2e + U_i)f(2e + U_i, t) \]
describes how an electron with kinetic energy \( \epsilon \) greater than the bandgap energy \( \nu_e \) uses this energy to excite a second electron from the valence to the conduction band. The three body recombination term obtained by detailed balance from the impact ionization model describes the process in which two electrons are lost from energy level \( \epsilon \) - one of them drops to the valence band, while the other is promoted to a higher energy level with energy \( \epsilon + U_i; R_{rec}(e, t) = 2^{3/2} r(e) \theta(e/U_i - 1)f^3(e/U_i - 1, t) - 2^{3/2} r(2e + U_i)f^3(e, t) \). The coefficients \( v_i(e) \) and \( r(e) \) are the ionization and three-body recombination rates respectively and are detailed in [1]. The multiphoton (MPI) ionization rate creating the seed electrons for the avalanche ionization in the conduction band \( R_{mp}(e, t) \) is proportional to the laser intensity \( p(t) \times t^m \), where \( m \) is the minimum number of photons needed for ionization. An eight photon MPI term is used in our calculation which corresponds to 1053 nm wavelength of the laser field and the dielectric fused silica. Equation (1) is solved with the boundary conditions \( j(e, t) \bigg|_{e=0} = 0 \) and \( f(e, t) \bigg|_{e=0} = 0 \) restricting the electron diffusion into and out of the conduction band and limiting how high in the continuum electron distribution is allowed. The parameters for the numerical calculation are all listed in detail in [1] and [4]. The critical electron density for damage \( n_{cr} = 10^{17} \text{cm}^{-3} \) [1], [3] is defined as the density for which the electron plasma becomes reflective. At this electron density a damage spot which is a permanent modification of the material surface is observable at a location corresponding to the peak of the Gaussian spatial profile. We use a one dimensional finite element method for solving the Fokker-Planck kinetic equation with its pertaining boundary conditions, and for each pulse duration \( \tau_p \) adjust the maximum laser intensity \( I_o \) in such a way that the laser fluence, defined as \( H = \int_{-\infty}^{+\infty} I(t)dt \approx I_o \tau_p \), produces the critical electron density \( n_{cr} \). The value of the laser fluence which has produced \( n_{cr} \) for a given pulse duration is defined as the optical damage threshold fluence.

The time dependence of the electron number density is plotted along with the Gaussian pulse shape in figure 1 for a 100 fs pulse. The laser intensity necessary to achieve critical electron density for a 100 fs pulse is \( I_o = 13 \text{TW/cm}^2 \). For \( \tau = 100 \text{fs} \) we plot the scaled electron distribution function as a function of the scaled energy for different scaled times (figure 2). The laser pulse is centred at scaled time \( t = 0 \) and the run starts at \( t = -1 \). For early times (\( t < 0 \)) the initial electron distribution function is established by the MPI term. When the avalanche starts, the distribution function increases in magnitude and its shape changes for different times. When the laser intensity decreases and the contributions from the intensity dependent diffusion, joule heating and MPI terms become less significant the electron distribution concentrates closer to \( \epsilon = 0 \) due to the fact that this is the regime in which the energy loss, thermal diffusion and recombination terms dominate over the intensity dependent terms and with the impact ionization eventually reduce the distribution function to zero [1].

We study the sensitivity of the electron density and threshold fluence to the input rates of the Fokker-Planck equation (1) for two pulse durations so that suitable materials may be developed that have the appropriate damage threshold for certain applications. First we find the intensities \( I_o \) producing \( n_{cr} \) for the chosen pulse duration using the standard parameter set (used originally to solve the equation).
Figure 1. Electron density $n_e$ (scaled in $n_{cr}$) as a function of the time (scaled in $\tau_e$) for laser pulse with $\tau_e=100fs$ and the laser intensity profile for reference.

Figure 2. Energy (scaled in $U_j$) spectra of the electron distribution function $f$ (scaled in $n_e$ and $U_j$) for different scaled time for laser pulse with $\tau_e=100fs$.

Then for the same intensity and for the given pulse duration we vary the rates one at a time and show what effect that has on the electron density. To obtain the change in threshold fluence we adjust the laser intensity for each varied rate to obtain scaled $n_{cr}=1$. The diffusion term $D$ spreads the electron distribution function to lower and higher energies, so that larger $D$ should spread the electron density faster. The more rapid electron diffusion to higher energy levels enhances the effect of the avalanche and hence increases the electron number density and vice versa. As seen from figures 3 and 4 the results we get from the $D$ term sensitivity test are consistent with our expectations. Comparing

Figure 3. The variation of electron density in the sensitivity test for the diffusion term $D$ for a laser pulse with $\tau_e=25fs$, wavelength $\lambda=1053nm$ and threshold intensity $I_0=22.7TW/cm^2$.

Figure 4. The variation of electron density in the sensitivity test for the diffusion term $D$ for a laser pulse with $\tau_e=300fs$, wavelength $\lambda=1053nm$ and threshold intensity $I_0=6.95TW/cm^2$. 
figures 3 and 4 also suggests that for $\tau_2 = 300\,\text{fs}$ the diffusion term $D$ is more significant for the development of the avalanche than for $\tau_2 = 25\,\text{fs}$. This is also confirmed by the change in the threshold fluence $H_{th}$ which is 10% for $\tau_2 = 25\,\text{fs}$ and 26% for $\tau_2 = 300\,\text{fs}$.

The double pulse calculation is carried out in part to show how fast the recombination terms will deplete the electron density created by the first laser pulse when its intensity is below $I_{cr}$, and how this will affect the second pulse intensity needed to damage the dielectric for different delays between the two pulses. As explained in detail in [1], [4] we introduce different recombination factors $\eta_1=1$ and $\eta_2=10$. Figure 5 illustrates that when the first pulse intensity $I_1$ is 85% of the single pulse threshold intensity ($I_0$) the electron density decays before the second pulse irradiates the sample for all three (scaled in $\tau_L$) delay times. Figure 6 shows how the ratio $1-r_2=1-I_2/I_0$ depends on the ratio $r_2 = I_2/I_1$ for a given delay time, where $I_0$, $I_1$ and $I_2$ are the single pulse threshold intensity, the first pulse intensity and the second pulse intensity respectively. The ratio $1-r_2$ shows by how much the intensity needed for damage by a single pulse is reduced when a double pulse is used. It is seen from figure 6 that for $r_2 < 1$ we can get 50%-70% reduction in the single pulse damage fluence by using double pulse for delay time of 60/fs and factor $\eta_1=1$. This of course a theoretical prediction and we are aware of the fact that such a double pulse experiment is difficult to realize in practice.

![Figure 5](image1.png)  
**Figure 5.** Electron density (scaled in $n_{cr}$) as a function of the pulse delay time (scaled in $\tau_L$) for laser pulses with $\tau_2=25\,\text{fs}$, $I_0=77\,\text{TW/cm}^2$, $\tau=25\,\text{fs}$, $I_1=77\,\text{TW/cm}^2$ and $\eta_1=1$. $I_1=65.45\,\text{TW/cm}^2$, $I_2=13\,\text{TW/cm}^2$ and $\eta_2=10$.

![Figure 6](image2.png)  
**Figure 6.** The ratio $1-r_2=1-I_2/I_0$ as a function of the ratio $r_2 = I_2/I_1$ for laser pulses with $\tau_2=25\,\text{fs}$, $I_0=77\,\text{TW/cm}^2$, $\tau=25\,\text{fs}$, $I_1=77\,\text{TW/cm}^2$ and $\eta_1=1$. $I_1=65.45\,\text{TW/cm}^2$, $I_2=13\,\text{TW/cm}^2$ and $\eta_2=10$.

4. Laser interaction with semiconductors

Next we consider a pulsed laser radiation incident on an undoped bulk GaAs semiconductor material at finite lattice temperature and with temperature dependent bandgap energy $E_G$. We start with the total Hamiltonian of the electron-phonon system exposed to a pulsed laser field in second quantization [5]. Initially using Fermi’s Golden Rule we obtain a dynamic equation by including only the electron phonon interaction within the first order perturbation theory. Limiting ourselves to the diffusive limit
\[\hbar \omega_q \ll E^* \epsilon, \quad \text{where} \quad E^* \epsilon = \hbar^2 k^2 / 2 m^* \epsilon \quad \text{and} \quad \hbar \omega_q \] are the free electron kinetic energy in the conduction band with effective mass \( m^* \epsilon \) and the phonon energy respectively, we expand the electron distribution function and get the Fokker-Planck equation for the thermal motion of electrons in energy space. We find that the intraband transitions of conduction electrons cannot optically respond to the spatially uniform laser field in a direct way due to the impossibility of conserving energy and momentum simultaneously during the absorption of a photon by an electron, and we include the effect of the laser field treated classically in the equation using the following considerations:\( (a) \) after the laser field is applied there is no net momentum change of electrons within an interval of collision time \( \tau_p \) (many time periods of the oscillating laser field) over which two successive collisions with phonons occur, and this is different from the case with a \( dc \) electric field; \( (b) \) electrons will gain net power from the laser field. In the presence of the laser field given by \( E_\lambda(t) = E_{\text{lo}} \cos(\Omega t) \), where \( E_{\text{lo}} \) and \( \Omega = 2\pi / \lambda \) are the amplitude and angular frequency, an energy-space drift exists which is described by:

\[
\frac{\partial f^*}{\partial t} = \lim_{\Delta t \to 0} \frac{1}{\Delta t} \left[ f^*(E^* \epsilon, E^* \epsilon, \Omega t) - f^*(E^* \epsilon, E^* \epsilon, \Omega t) \right] - \left( \frac{\partial E^* \epsilon}{\partial E^* \epsilon} \right) f^*(E^* \epsilon, E^* \epsilon, \Omega t).
\]

Since the time period of the laser field \( 2\pi / \Omega \) is much shorter than \( \tau_p \), a time average is taken of the Fokker-Planck type equation over many periods of the laser field, an averaged fluctuation which represents the average energy gained by electrons by absorbing the laser power \( E_{\text{lo}} \ll \Delta E^* \epsilon \) is introduced and for \( E_{\text{lo}} \ll E^* \epsilon \) the expansion \( f^*(E^* \epsilon, \Delta E^* \epsilon) \approx f^*(E^* \epsilon) - E_{\text{lo}} \frac{\partial f^*}{\partial E^* \epsilon} \) is used. Finally we obtain the generalized Fokker-Planck type equation

\[
\frac{\partial f^*}{\partial t} + V^* \frac{\partial E^* \epsilon}{\partial E^* \epsilon} - D^* \frac{\partial^2 f^*}{\partial (E^* \epsilon)^2} = A_{ij} f^* + S_i.
\]

Its coefficients and the parameters used in the numerical calculations for a bulk GaAs semiconductor sample are detailed in [5], [6]. In this equation we have included the coupling of the energy drift of electrons in the presence of the laser field to the phonon-assisted intraband electron transitions which gives rise to an additional anti-diffusion current and have given analytical expressions for all the source terms up to second order perturbation theory, including stimulated interband electron transitions due to a single-photon absorption, impact ionization due to Coulomb interaction between electrons and holes, and non-radiative recombination due to a phonon-mediated interaction. The new anti-diffusion term describes a process in which the spontaneous phonon emission from electrons decreases when the slope of the distribution function is positive and increases when its slope is negative, i.e. it is opposite to the diffusion process. When an incident light field is absorbed by an undoped bulk semiconductor (dielectric), the electrons in the completely filled valence band are excited to the conduction band and the free electron density increases with the intensity of the incident light. The electrons in the conduction band form electron plasma when the density is high. When the condition \( \Omega^2 \epsilon \leq e^2 n^*/\epsilon_0 \epsilon, \) with \( \epsilon, \ epsilon_0 \) and \( \epsilon \) being the electron charge and the dielectric and relative static dielectric constants respectively, is satisfied, the incident light will be totally reflected by the surface of the semiconductor and the electron density reaches its critical (maximum) value \( n_{\text{cr}} \). In this case the semiconductor (dielectric) is optically damaged which is measured by observing how the semiconductor device degrades in performance (ablation-actual removal of a thin layer of material for dielectrics). In the case of GaAs for \( \hbar \Omega \epsilon = E_{\text{lo}} = 1.42eV \), we get \( n_{\text{cr}} = 1.3 \times 10^{20} \text{cm}^{-3} \). An undoped semiconductor behaves like an insulator but becomes a good conductor at room temperature (\( T=300K \)). The thermally excited conduction electron density in an undoped semiconductor in equilibrium is \( n = \sqrt{n_v n_h} \epsilon^{3/2} \), where \( n_v = 2(m^* \epsilon k_B T / 2\pi \hbar^2)^{3/2} \) are the state densities of holes and electrons respectively. When the
temperature $T$ is equal to the lattice melting temperature $T_m$, $n_f$ reaches its maximum value $n_{elec-cr}$ or electrical limit and in this case the semiconductor is electrically damaged. Since for GaAs $T_m=1512 K$, we get $n_{elec-cr}=8.7 \times 10^{17} \text{cm}^{-3}$. The atoms in most semiconductors, such as Si, Ge and GaAs, are chemically connected by covalent bonds with cohesive energy $E_{coh} \cong E_G$ in a crystal. When an undoped semiconductor is exposed to incident light field, the statistically averaged kinetic energy per electron, which is proportional to electron temperature $T_e$, is $<E^*_{e}> = \left[ \frac{1}{2m_e} \int_0^{E_G} E f^e_{cr} dE e^*_{cr} \right]$ When the electron distribution function $f^e_{cr}$ peaks at increasingly higher energies, $<E^*_{e}>$ also increases although the conduction electron density can be very small at this time. When $<E^*_{e}> = E_G$, there is an instability in the chemical bonds of the semiconductor and the semiconductor is structurally damaged by the laser field for semiconducting-material uses. For a peak laser intensity $I_L = 8 \times 10^9 \text{W/m}^2$ and laser pulse length $\tau_L = 1 \text{ps}$, we find that the calculated conduction electron density $n_e$ satisfies the inequality $8.7 \times 10^{17} \text{cm}^{-3} = n_{elec-cr} < n_e < n_{opt} = 1.3 \times 10^{23} \text{cm}^{-3}$. This means that the semiconductor GaAs has undergone electrical damage but is not optically damaged. Our calculations also show that the average kinetic energy per electron $<E^*_{e}>$ is always less than $E_G$, which implies that the sample GaAs is structurally stable. The role played by the different terms of the equation (thermal emission, recombination, diffusion, etc.) in determining the dynamics of the electron distribution function in energy space has been analyzed below. We see in figure 7 that when $t/\tau_L = -0.5$ a feature is developed in the electron distribution function $f^e_{cr}$ on the low energy side of the peak as a result of the competition between the upward anti-diffusion and the downward thermal diffusion (large at low energies and high temperatures) when $E_m \approx k_b T$.

**Figure 7.** Comparison of energy spectra of $f^e_{cr}$ at $t/\tau_L = -0.5$ for including and excluding the anti-diffusion current.

**Figure 8.** Comparison of energy spectra of $f^e_{cr}$ at scaled time of the laser pulse for including and excluding thermal emission.

Thermal emission is the main energy loss mechanism, so the electrons from high energy states in the conduction band are transferred to low-energy states through emission of spontaneous phonons shifting the peak down. Since the transfer of electrons to the bandedge enhances the upward
anti-diffusion current, a pronounced feature appears at the low energy side of the peak as can be seen in figure 8. Since the single-photon absorption is a resonant process, the peak of the distribution function diffusion current, a pronounced feature appears at the low energy side of the peak as can be seen in figure 8. Since the single-photon absorption is a resonant process, the peak of the distribution function appears initially at the energy determined by the laser frequency detuning away from the band-edge. At early times the peak position is solely determined by the detuning.

As seen in figure 9 the peak position at later times is determined by a balance among the different diffusion currents, thermal emission and joule heating so the peak position becomes independent of the detuning. Recombination results from phonon-mediated electron-hole interaction. From the one-photon excitation we expect to find a spike-like feature at the bandedge, but in figure 10 we find that the recombination term reduces the occupation of the electrons at the bandedge hence suppresses the expected spike-like feature to a kink.

Figure 9. Comparison of energy spectra of \( f^e_k \) at \( t/\tau_L = -0.5 \) for a larger (solid curve) or smaller (dashed curve) laser frequency detuning.

Figure 10. Comparison of energy spectra of \( f^e_k \) at \( t/\tau_L = 0.1 \) for including and excluding non-radiative recombination.

5. Laser interaction with quantum wells

We show that an anisotropic and non-linear renormalized electron-phonon interaction arises in GaAs/AlGaAs quantum wells in the presence of a normally incident, intense mid-IR pulsed laser field. The pulsed laser field does not directly couple to intersubband electron transitions due to the plane-polarized field and small photon energy compared to the energy separation between the two lowest subbands in the quantum wells. The laser field increases the electron kinetic energy, and the increased electron energy can be exchanged with phonons or used to promote the electrons to higher subbands. These processes significantly change the population in both subbands. We consider a uniformly doped quantum well with electron confinement in the \( z \) (growth) direction. For a uniform electromagnetic field polarized in the quantum-well plane perpendicular to \( z \) we assume a vector potential \( \vec{A}(t) = [\phi(t),0,0,0] \), where \( \phi(t) = \sin(\Omega_L t) \), and \( \Omega_L \) is the angular frequency of the laser field and \( \phi(t) \) is a step function, as a result of the rotational invariance of the system in plane. For the single-particle Hamiltonian of the system in the presence of the given above vector potential, using the time dependent Schrödinger equation we get the total wave function for conduction electrons. Using the electron wave function and first order time-dependent perturbation theory we find the transition rate between initial and final electron states [7]. The intrasubband electron transitions cannot directly respond to the uniform laser field due to the violation of momentum conservation. However, we find...
from the expression for the transition rate between initial and final electron states given in [7], [8] that with the help of phonon scattering incident photons can be absorbed by intrasubband (n = n) (where n is the index of the quantum well subbands) electron transitions. The intersubband electron transitions can not directly respond to the plane-polarized laser field through dipole coupling due to selection rules. Although we find from expression for the transition rate between initial and final electron states that the intersubband optical-phonon scattering of electrons (M = 0) becomes possible due to renormalization of the electron-phonon interaction in the quantum well by the laser field where M is the order of the renormalizing Bessel function. In the above expression \( \tilde{\mathbf{k}}_r = (\tilde{k}_r, \tilde{k}_r) \) and \( \tilde{\mathbf{q}}_r = (\tilde{q}_r, \tilde{q}_r) \) are the two dimensional electron and phonon wave vectors in the quantum well plane respectively, \( E_n(k_r) = E_{k_r} + E'_n \) with \( E_{k_r} = \hbar^2 k_r^2 / 2m_r \) and \( E'_n \) the edge of the n\textsuperscript{th} subband in the quantum well, and \( h\omega_{LO} \) is the energy of the longitudinal optical phonons. Finally for a given electron state, from the calculated transition rate we obtain the total scattering-in rate and scattering-out rate of conduction electrons, from which we arrive at the Boltzmann equation for the electron distribution function \( f_n(E_n^*, E_n, t) \) in each subband of the considered quantum well. The parameters of the numerical calculation for the two selected samples are given in detail in [7].

The pulsed laser is chosen to have a Gaussian profile in time with a pulse duration of \( 1\text{ps} \). In all the figures we use the notation \( h\Omega'_{21} = E'_n - E'_n \) to represent the energy difference of the two subband edges and we provide the value of the wavelength used by giving the numerical value of the photon energy \( h\Omega_z \) [7]. In the figures shown below the magnitude of the photon energy is \( h\Omega_z = 56.5meV \).

Figure 11 shows the difference between the electron distribution functions of sample 1 (see [7]) with/without the photon-assisted process for non-resonant excitations at the moment of peak pulse \( t = t_p \). For the second subband (n = 2) the change of electron distribution function (dashed curve) exhibits two upward steps at \( E_{k_r} = h(\Omega_z + \omega_{LO} - \Omega'_{21}) = 10.1meV \) and \( E_{k_r} = h(\Omega_z - \omega_{LO}) = 20.2meV \). The higher energy step (upward arrow) reflects the strong photon

\[ h(\Omega_z + \omega_{LO} - \Omega'_{21}) \]

\[ h(\Omega_z - \omega_{LO}) \]

\[ h\Omega_z = 56.5meV \]

\[ T = 200K \]

\[ n = 1 \]

\[ n = 2 \]

\[ h\Omega_z = 56.5meV \]

\[ T = 150K \]

\[ n = 1 \]

\[ n = 2 \]

**Figure 11.** Difference between the electron distribution functions of sample 1 with/without the photon-assisted process at \( t = t_p \) for \( f_1(E^*_{n} + E_{k_r}, t) \) (solid curve) and \( f_1(E^*_{n} + E_{k_r}, t) \) (dashed curve) with set of parameters given in [7].

**Figure 12.** Difference between the electron distribution functions of sample 1 with/without the photon-assisted process at \( t = t_p \) for \( f_1(E^*_{n} + E_{k_r}, t) \) (solid curve) and \( f_1(E^*_{n} + E_{k_r}, t) \) (dashed curve) with set of parameters given in [7] and \( T = 150K \).
absorption induced phonon emission by intrasubband transitions of electrons within the second subband, while the lower-energy step (downward arrow) indicates the effect of the photon-absorption induced phonon absorption by intersubband transitions of electrons from the first to the second subband. For the first subband \( (n=1) \), from the difference between the electron distributions with/without photon assistance (solid curve) we find only one upward step at \( E_n = h(\Omega - \omega) \). This higher energy step also comes from the strong photon-absorption induced phonon emission by intrasubband transitions of electrons within the first subband.

Figure 12 shows the same as depicted in figure 11 except for a different lattice temperature \( T = 150 K \). Compared with figure 11, the photon-absorption-induced phonon absorption by intersubband transitions of electrons from the first to the second subband (downward arrow) at \( T = 150 K \) is reduced, as is the photon-absorption-induced phonon emission within the second subband (upward arrow) for \( (n = 2) \) (dashed curve). For \( (n = 1) \) (solid curve) we find only insignificant change compared with figure 11. We present more figures similar to 11 and 12 in reference [7] in which we vary the photon energy, the laser field strength, and the quantum well width. We find out that the decrease in the amplitude of the laser field \( \varepsilon_0 \) only slightly reduces the effects of the photon-absorption-induced phonon emission and absorption. Since the factor renormalizing the electron-phonon interaction in the Boltzmann equation is \( R \propto \varepsilon_0/\Omega^2 \), we expect a smaller reduction in photon-assisted effects when \( \varepsilon_0 \) decreases in comparison to the case when the photon energy \( h\Omega \) increases. The increase in photon energy weakens the effect of photon absorption induced phonon emission and absorption. The increase of \( h\Omega \) also pushes all steps in the distribution functions (upward and downward arrows) to higher energies.

6. Conclusions
In summary we have explored theoretically the effects of the interaction of ultra-short high power lasers with dielectrics, semiconductors and semiconductor quantum wells using the kinetic Fokker-Planck and Boltzmann equations. The energy spectra of the electron distribution function at different times relative to a laser pulse and the conduction electron density as a function of time were obtained to explain the role played by the different terms in the equations corresponding to different physical processes taking place during the interaction, such as thermal and field dependent diffusion, anti-diffusion, Joule heating, recombination, thermal emission, single photon and multiphoton ionization, impact ionization in the case of dielectrics and semiconductors, and the occurrence of a forbidden photon absorption for a plane polarized laser field with photon energy smaller than the subband width \( (h\Omega \ll h\Omega_n) \) and a “forbidden” intersubband phonon-scattering of electrons for phonon energy smaller than the subband width \( (h\omega_n \ll h\Omega) \) in the case of semiconductor quantum wells.

The results presented are not limited to the samples used in the calculations, and the described theory can be used or developed further for the needs of laser microprocessing and quantum electronics.

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