Ultrafast high-field transport is studied in $n$-type GaAs with ultrashort terahertz (THz) pulses of an electric field amplitude of up to 300 kV/cm. At lattice temperatures between $T=300$ and 80 K, we observe coherent ballistic transport of electrons over a major part of the first Brillouin zone. At $T=300$ K, ballistic transport occurs at a constant electron density whereas at lower temperatures, the THz pulses generate additional electron-hole pairs by field-induced tunneling between valence and conduction bands. We show that the ultrashort decoherence time of superpositions of valence- and conduction-band states plays a crucial role for the efficiency of the tunneling process. The extremely fast interband decoherence at room temperature results in a negligible tunneling rate.

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I. INTRODUCTION

In the periodic potential of a crystal, an electron under the action of an electric field performs a ballistic motion if scattering processes have a negligible influence. In this case, the change in electron wave vector $\vec{k}$ with time is solely determined by the electric field acting on the electron,\(^1,\(^2\)

$$\frac{d\vec{k}}{dt} = -e\vec{E}(t).$$

(1)

From the time-dependent wave vector, the momentum electron velocity is obtained via the gradient of the relevant band in the band structure (Fig. 1),

$$\vec{v}(\vec{k}) = \frac{1}{\hbar} \nabla_{\vec{k}} \mathcal{E}(\vec{k}).$$

(2)

Under most circumstances, the prerequisite for Eqs. (1) and (2), namely, that scattering is unimportant, is not fulfilled. If scattering is important, we typically find drift transport, described by the Drude model\(^6\) and concepts such as electron mobility.

Very recently,\(^7\) we have demonstrated that electrons in bulk $n$-type GaAs at room temperature ($T=300$ K) undergo coherent ballistic transport in the high electric field of several 100 kV/cm of short terahertz (THz) pulses. The THz field radiated by the accelerated electrons demonstrates that the carriers are coherently driven into the range of the conduction band where the effective mass becomes negative. A quantitative analysis shows that the coherent electron motion covers approximately half of the first Brillouin zone of bulk GaAs. Such results clearly demonstrate negligible friction due to incoherent scattering processes, in particular, polar optical-phonon scattering,\(^8,\(^9\) on the few hundred femtosecond time scale of the experiments. The observed behavior is fully reproduced by quantum-kinetic transport calculations.

In this paper, we present an extension of the previous work\(^7\) toward lower lattice temperatures ($T=200$ K and $T=80$ K). The results presented in the following clearly demonstrate coherent carrier transport at such lower temperatures. In contrast to room temperature, the number of electrons undergoing such motion is substantially enhanced and exceeds the number of electrons present by doping. This shows that the transient THz field promotes a substantial number of electrons from the valence into the conduction band, although the THz photon energy is much less than the band gap, $\hbar\omega_{\text{THz}} < E_G$. Field-induced tunneling from the valence to the conduction band represents the main mechanism enhancing the number of free carriers. The tunneling rate depends sensitively on the decoherence rate between valence- and conduction-band states. The temperature depen-
To study high-field transport on ultrafast time scales, we need both high electric driving fields and a method to determine the electron velocity as a function of time.

In our experiment (Fig. 2), high electric fields are applied to the sample in the form of strong THz pulses. The time dependence of the electric field acting on the sample is measured by electro-optic sampling.\textsuperscript{10–13} The detection scheme is used to measure the electric field emitted by the electrons in the sample, which is in our geometry proportional to the current density

\begin{equation}
    j(t) = -en(t)v(t),
\end{equation}

where \(n(t)\) and \(v(t)\) represent the time-dependent density and velocity of electrons.

The sample investigated was grown by molecular-beam epitaxy and consists of a 500 nm thick layer of \(n\)-type GaAs with a donor (Si) concentration of \(N_D=2 \times 10^{16} \text{ cm}^{-3}\) clad between two 300 nm thick \(Al_{0.4}Ga_{0.6}As\) layers (the same sample was studied in Ref. 7). In the relevant area, the substrate is etched away\textsuperscript{14} so that the experiment is performed on a free-standing GaAs film.

THz pulses with a high electric field amplitude are generated by optical rectification of 800 nm pulses from a femtosecond Ti:sapphire oscillator-amplifier system. An acoustooptic pulse shaper\textsuperscript{15,16} between oscillator and amplifier allows to tune both spectral amplitude and spectral phase to obtain a few-cycle THz pulse with a center frequency of 2 THz, which excites the sample placed in the focus of a pair of parabolic mirrors. The electric field of the transmitted THz pulse is measured via electro-optic sampling in a thin ZnTe crystal. The entire optical path of the THz beam is placed in vacuum.\textsuperscript{7–19} For the low-temperature measurements, the sample is mounted on the cold finger of a helium-flow cryostat, the latter being also placed in the vacuum chamber.

The electron current density,

\begin{equation}
    j(t) = \frac{-2E_{\text{em}}(t)}{Z_0d},
\end{equation}

(see Appendix) in the sample is proportional to the coherently emitted field

\begin{equation}
    E_{\text{em}}(t) = E_{\text{in}}(t) - E_{\text{at}}(t),
\end{equation}

which is given by the difference of \(E_{\text{en}}(t)\), the field transmitted through the sample, and \(E_{\text{at}}(t)\), the field incident on the sample \((Z_0=377 \text{ \Omega, the impedance of free space})\). As the thickness of our sample \(d=500 \text{ nm}\) is much less than the THz wavelength \(\lambda=150 \text{ \mu m}\), all electrons in the sample experience the same driving field \(E_{\text{loc}}(t)\), which is identical to \(E_{\text{en}}(t)\).\textsuperscript{18,20} The two transients \(E_{\text{en}}(t)\) and \(E_{\text{at}}(t)\) are measured in separate scans, \(E_{\text{af}}(t)\) with the sample in place and \(E_{\text{at}}(t)\) without the sample. Since the sample is very thin, the delay introduced by its refractive index has a very small value of 8 fs. This delay is taken care of when calculating \(E_{\text{em}}(t)\), Eq. (5).

\section*{III. RESULTS}

In Fig. 3, we show the incident and the emitted electric field transients measured at sample temperatures of 200 and 80 K for different amplitudes of the incident field. In all cases, one finds that in the beginning of the pulse \(E_{\text{en}}(t)\) and \(E_{\text{at}}(t)\) are out of phase, signifying absorption of the THz pulse in the sample. At later times, \(E_{\text{em}}(t)\) and \(E_{\text{at}}(t)\) are in phase, signifying THz emission.

Of particular interest is the comparison of the present low-temperature results with the results at room temperature, which have been discussed in detail in Ref. 7. At \(T=300 \text{ K}\), the amplitude of the emitted electric field is limited to about 7 kV/cm. The analysis in Ref. 7 shows that this amplitude agrees with the maximum emitted field possible for the electron density present by doping and the maximum electron velocity reached in ballistic transport (see Fig. 3 in Ref. 7). In contrast, at an incident amplitude of 300 kV/cm, the emitted field amplitude reaches 35 kV/cm at \(T=200 \text{ K}\) [Fig. 3(b)] and more than 50 kV/cm at \(T=80 \text{ K}\) [Fig. 3(d)]. In Fig. 4, we show \(E_{\text{em}}(t)\) at \(T=80 \text{ K}\) for an incident amplitude of 300 kV/cm on an extended scale and compare it with the results at room temperature under otherwise identical conditions.

According to Eqs. (3) and (4), the emitted field is proportional to the electron density times electron velocity. The velocity is determined by the band structure [Eq. (2)], which
changes only slightly as a function of lattice temperature. We thus conclude that the much higher amplitude of the field emitted at lower temperatures is due to an increase in the density of carriers contributing to the emission. Such an increase in electron density requires the promotion of electrons from the valence into the conduction band, i.e., an excitation energy equal to at least the band-gap energy $E_g = 1.5$ eV. To find out whether the THz pulse can provide the required energy density to the sample, we calculate the transiently absorbed energy\(^\text{18}\) normalized to the density $N_D$ of electrons present by doping,

$$W_{\text{abs}}(t) = \frac{1}{N_D} \int_{-\infty}^{t} j(t')E_i(t')dt'$$

$$= \frac{-2}{dN_D\mu_e} \int_{-\infty}^{t} E_{\text{em}}(t')E_i(t')dt'. \quad (6)$$

In Fig. 4(b), $W_{\text{abs}}$ for $T=80$ K is plotted as a function of time (solid line). In the following, we consider two particular values of $W_{\text{abs}}(t)$, the maximum $W_{\text{max}}$ and the value $W_{\text{tr}}$, the irreversibly absorbed energy after the end of the THz pulse.

**IV. MECHANISM FOR THZ-INDUCED INTERBAND TRANSITIONS**

There are two possible mechanisms for the transfer of electrons from the valence into the conduction band, impact...
ionization and tunneling. An important issue for distinguishing between the two possibilities is whether the mechanism explains the experimentally observed temperature dependence, i.e., a negligible increase in electron density in the conduction band at room temperature and a strong increase at lower temperatures ($T=80$ and 200 K).

In impact ionization, an electron loses its kinetic energy to generate an electron-hole pair. This is only possible if the electron has a kinetic energy larger than the band gap. Theoretical descriptions of impact ionization predict ionization rates that depend only on the electron kinetic energy and on the band structure. Thus, the lattice temperature influences the ionization rate only indirectly, either by changing the band structure or by preventing the electron to reach a sufficient kinetic energy. Both effects cannot explain our results. On the one hand, the band structure changes only if only the electrons present by doping participate, i.e., if no tunneling occurs.

FIG. 5. (Color online) Maximum energy $W_{\text{max}}$ (open symbols) and irreversibly absorbed energy $W_{\text{irr}}$ (closed symbols) as a function of the incident electric field amplitude, calculated according to Eq. (6), for temperatures of 80, 200, and 300 K. For $T=80$ K and for $T=200$ K, the dashed lines are fits to $W_{\text{max}}=\text{const } E_{\text{in}}^2$. The hatched area shows the possible values for $W_{\text{max}}$ if only the electrons present by doping participate, i.e., if no tunneling occurs.

part of the pulse, impact ionization can at most account for the generation of 0.1 electron-hole pairs per initially present electron. Compared to the experimentally observed generation of ten electron-hole pairs this again shows that impact ionization is not able to explain our results.

Tunneling of electrons in high electric fields from the valence band into the conduction band is a well-known effect (Zener tunneling). Using the theoretical results for Zener tunneling in a static electric field, one finds a temperature-independent tunneling rate much lower than the rate needed to explain our results. For example, a tunneling rate close to zero is derived from Eq. (39) of Ref. 31 in a static electric field of 300 kV/cm. However, the model of Refs. 31–33 is not applicable to the present experiment as it neglects decoherence processes, i.e., the decay of quantum coherences between carrier wave functions of the initial and final states. As shown by Kazarinov and Suris (see also Ref. 36), decoherence strongly influences the tunneling rate $R$.

$$R \approx \frac{\tau}{1 + (\Delta \omega \tau)^2}. \quad (7)$$

In this equation, the decoherence is described by the constant decoherence time $\tau$ (we shall see below that actually $\tau$ will not be constant), $\Delta \omega$ is the detuning from resonance of the initial and final state. For resonant tunneling ($\Delta \omega=0$) Eq. (7) shows that decoherence decreases the tunneling rate (the highest $R$ is obtained for the longest $\tau$). For nonresonant tunneling, however, the tunneling rate will increase up to the point where the decoherence time $\tau$ is equal to the inverse of the detuning $\Delta \omega$. For even shorter decoherence times, the tunneling rate decreases again [see Fig. 6].

In our case, the detuning is approximately $E_{\text{in}}/h$, since the THz frequency is much smaller than the band gap. Thus, the highest tunneling rate is obtained for $\tau=0.5 \text{ fs}$. Under the reasonable assumption that $\tau$ decreases with increasing temperature, the decoherence times at the temperatures of our experiment have to be near the values marked in Fig. 6 to explain our results.

FIG. 6. (Color online) Population of the upper level (conduction band) after the end of the exciting THz pulse calculated for a two-level system as a function of the constant decoherence time $\tau$ (dots). The diamonds are estimates for the decoherence times at the temperatures of our experiment.
In this equation, $\tau_m$ is the momentum relaxation time, responsible, e.g., for low-field transport. In contrast to the model used in Eq. (7) and for Fig. 6, the decoherence rate in Eq. (8) is not constant in time since it depends on the time-dependent $\Delta x$. For a temperature of 80 K, one finds from the mobility $\tau_m=1000$ fs. At the maximum $\Delta x$, this results in the very short decoherence time of 0.1 fs. Thus, Eq. (8) accounts (i) for the very short decoherence times needed for high tunneling rates and (ii) for the temperature dependence of decoherence (apart from the direct proportionality of the decoherence rate on $T$, $1/\tau_m$ increases with $T$). On the other hand, Eq. (8) still allows for sharp features in linear interband absorption since there $\Delta x$ will be very small.

In contrast to the case considered in Refs. 34–36, Eq. (8) demonstrates a pronounced temperature dependence of $\tau$. Taking this fact into account, we performed calculations considering on the same footing both interband and intraband transitions under the influence of an electric field. The electronic band structure was included in the form of a pseudopotential\(^{42}\) calculation neglecting spin-orbit interaction with a finite set of local pseudopotentials. The general Bloch wave function of an electron with wave vector $\vec{k}$ in band $b$,

$$\psi_{b,\vec{k}}(\vec{r}, t) = e^{i\vec{G}\cdot\vec{r}} \sum_{\vec{G}} c_{b,\vec{k},\vec{G}}(t) e^{i\vec{G}\cdot\vec{r}} ,$$

is a coherent superposition of plane waves displaced by the reciprocal-lattice vectors $\vec{G}$. The coefficients $c_{b,\vec{k},\vec{G}}(t)$ are obtained from the following Hamiltonian, which includes the interaction with the external electric field via the vector potential $\vec{A}(t)=\int_0^t \vec{E}(t') dt'$ and the interaction with the periodic crystal potential:\(^{42}\)

$$[H^s(\vec{k}, \vec{A})]_{\vec{G}, \vec{G}'} = \frac{\delta_{\vec{G}, \vec{G}'}}{2m_0} \left[ h(\vec{k} + \vec{G}) - e\vec{A} \right]^2 + V(\vec{G} - \vec{G}').$$

The $V(\vec{G})$ are the coefficients of the Fourier expansion of the periodic potential and $m_0$ is the free-electron mass. By diagonalizing the Hamiltonian $H^s$, one can transform it into the $\vec{k}$- and $\vec{A}$-dependent eigenbasis $H^p$, $H^p(\vec{k}, \vec{A}) = S(\vec{k}, \vec{A}) H^s(\vec{k}, \vec{A}) S^{-1}(\vec{k}, \vec{A})$. For $\vec{A}=0$, the Hamiltonian yields the single-electron band structure, i.e., one gets the band energies $E_b(\vec{k})$ and the corresponding coefficients $c_{b,\vec{k},\vec{G}}(t)$ for the band $b$.

In our calculations, we work in the single-particle picture, i.e., we neglect all many-body interactions, among them electron-electron scattering,\(^{43}\) which could mix Bloch waves with different $\vec{k}$ vectors. Without this mixing, each electron can be labeled by its initial $\vec{k}_0$ and band index $b_0$ before THz excitation. To include decoherence according to Eq. (8), we calculate the electron-hole distance as the time integral of the difference of the electron and hole velocities. The velocities in turn are obtained from the band structure [Eq. (2)]. For simplicity, we use a single decoherence rate for all electron-hole coherences, using the average electron-hole distance.

In the initial state of our system, all valence-band states are filled and all conduction-band states are empty.\(^{44}\) The time evolution of the density matrix is obtained from the Liouville equation,
\[
\frac{dp}{dt} = \frac{1}{i\hbar} [H(t), \rho] - \Gamma(t, \rho). \tag{11}
\]

The last term in the equation above describes the decoherence. In principle, it is also necessary to consider spontaneous (incoherent) electron-hole recombination but this process can be neglected because of its low rate (typical lifetimes of electrons in the conduction band are in the nanosecond range). In solving Eq. (11), it is necessary to transform back and forth between the bases \(g\) and \(e\) mentioned above, using the transformation matrices \(S\) and \(S^{-1}\). The first term in Eq. (11) is calculated in the basis \(g\), the decoherence in the basis \(e\). For the numerical calculation, the size of the time steps has to be below the shortest \(\tau\) [Eq. (8)], i.e., at room temperature almost 0.01 fs, which makes the calculations very time consuming. From the calculation, one obtains the time dependence of the density matrix for the initial \(\vec{k}\). The resulting current density is obtained as the expectation value of the velocity operator times the electronic charge and the emitted electric field from Eq. (4). The main results are: (i) tunneling can only occur at such points in \(k\) space where the interband matrix element \(X_{b\rightarrow c}(\vec{k}) = \int \psi_{b, \vec{k}}^* \psi_{c, \vec{k}} dV\) is high. The band-structure calculation allows the determination of these interband matrix elements. It turns out that they are zero for the transition from the heavy hole to the conduction band so that only tunneling from the light hole to the conduction band contributes. The square of this matrix element \(X_{b\rightarrow c}^2(\vec{k})\) is shown in Fig. 1(c). One finds the maximum of this matrix element and thus of the tunneling rate at the center of the Brillouin zone. It decreases for larger wave vectors so that only electrons with wave vectors within a diameter of 1/10 of the Brillouin zone can tunnel into the conduction band. From the corresponding \(k\)-space volume and the population probability, we obtain a total electron density of \(3 \times 10^{17} \text{ cm}^{-3}\) in the conduction band, which agrees well with the observed value of \(2 \times 10^{17} \text{ cm}^{-3}\).

For ballistic transport, the electron wave vector is proportional to the vector potential. Thus, one expects interband tunneling near the zeros of the vector potential. An additional requirement for interband tunneling is a high electric field. The effect of these two competing requirements is shown in Fig. 8(a). We see that the increase in the conduction band population is especially high if a high electric field and a zero of the vector potential \(A\) coincide (at \(t=1.9\ \text{ ps}\)). In contrast, for impact ionization the highest increase in the conduction-band population is expected near the extrema of the vector potential.

(ii) As mentioned already, decoherence plays an important role. Without decoherence, any tunneling is completely reversible [curve for \(\gamma_m=0\) in Fig. 8(b)]. After the THz pulse, no electrons remain in the conduction band. For higher decoherence, we find an increasing conduction-band population after the end of the pulse, reaching \(4 \times 10^{-3}\) for \(\gamma_m=(1 \text{ ps})^{-1}\), the value relevant for our experiment at low temperature. The corresponding population probability at room temperature is \(5 \times 10^{-4}\), nearly a factor of 10 less.

(iii) Half of the additional current is carried by electrons in the conduction band, the other half by light holes. As can be seen in Fig. 1(b), the velocities in these bands are of comparable magnitude.

The theoretical curves in Fig. 4 (dashed lines) have been calculated with the model described above. While the time dependence of the absorbed energy \(W_{abs}(t)\) agrees quite well with the theoretical curve [Fig. 4(b)], the calculated emitted electric field shows deviations, in particular in the beginning of the pulse [Fig. 4(a)]. Nevertheless, for such a simple model, which neglects all many-body effects such as electron-electron scattering and the formation of excitons, the agreement is satisfactory.

V. CONCLUSIONS

In conclusion, ultrafast high-field THz transients induce ballistic electron transport in bulk GaAs. At room temperature, it is possible for the electrons to traverse half the Brillouin zone for electric field amplitudes of 300 kV/cm. At lower temperatures (80 and 200 K), we observe additional THz-induced electron-hole pair generation, which leads to an increase in the density of free carriers by a factor of 10. The mechanism of this electron-hole pair generation is field-induced tunneling from the valence into the conduction band. The tunneling rate depends strongly on the decoherence of the coherent superposition of electron and hole states. Since the strong electric field accelerates electrons and holes in opposite directions, we get a highly nonclassical wave function, which decoheres very fast because of its coupling to the environment.
APPENDIX: MEASUREMENT OF VELOCITY VERSUS ACCELERATION OF ELECTRONS ON ULTRAFAST TIME SCALES

The electromagnetic field radiated by a single charged particle in motion is essentially proportional to its acceleration, see, e.g., Eq. (14.18) in Ref. 46. In our present experiment and in previous experiments,47–49 however, we do not deal with a single moving particle but rather an extended sheet of simultaneously moving particles. As a first approximation, we assume an infinite sheet of particles all moving with the same velocity. Because of interference between the individual contributions of all particles, it turns out that the electric field at the sample19,20,50 is proportional to the velocity of the moving charges or, equivalently, to the current density \( j \) in the sample (see Ref. 19 for a direct derivation using Maxwell’s equations).

In the experiment, one does not directly measure \( E_{\text{sample}} \) but measures the electric field \( E_{\text{det}} \) at the electro-optic crystal. The relationship between \( E_{\text{det}} \) and \( E_{\text{sample}} \) depends on the optics used for the transfer of the electric field from the sample to the electro-optic crystal. We use a first parabolic mirror to collimate the radiation emitted from the sample and a second parabolic mirror to focus the radiation onto the electro-optic crystal [Fig. 9(a)]. Under the assumption of focusing mirrors collecting the entire solid angle, we get in this case \( E_{\text{det}} = -E_{\text{sample}} \). Thus, the detected field is proportional to the electron velocity.

We now consider how deviations of the actual experimental conditions from the idealized assumptions influence the results. Using the formalism of Ref. 50, it is possible to calculate the electric field at the electro-optic crystal generated by a Gaussian current distribution on the sample. It turns out that the on-axis field \( (\rho=0) \) is proportional to the on-axis current density down to spot sizes of about one THz wavelength and, for the mirror sizes used (diameter 25.4 mm and parent focal length of 12.7 mm), in the frequency range \( \approx 1 \) THz. This frequency range is perfectly adequate for our experiments since the incident field has negligible spectral density below 1 THz.

On the electro-optic crystal, the spot size of the probe beam is much smaller than the spot size of the THz beam. Therefore, we measure the on-axis THz field, as required by the analysis described. As a result, the signal measured in our setup is proportional to the electron velocity.

In contrast to our experiment, the setups of, e.g., Refs. 47–49 [Fig. 9(b)] use a single focusing element that images a large area of the sample onto a small spot on the electro-optic crystal. In the frequency domain, using Gaussian optics, one finds in this case \( E_{\text{det}}(\omega) \propto \omega E_{\text{sample}}(\omega) \). The reason for this is that the size of the focal spot is proportional to the wavelength and thus inversely proportional to the frequency. In the time domain, this leads to \( E_{\text{det}} \propto \partial E_{\text{sample}}/\partial t \). With \( E_{\text{sample}} \propto v \) one thus gets in this case that the signal is proportional to the acceleration of the particles.

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43 The justification for neglecting electron-electron scattering in the calculation is that electron-electron scattering conserves total energy and total momentum of the electron system. Thus, the influence of electron-electron scattering is only indirect, leading to a redistribution within the electron system, which may in turn influence the dynamics. Furthermore, even at the highest electron density the typical electric field an electron generates at the position of another electron is only 5 kV/cm, much less than the applied electric fields.
44 Compared to the total number of conduction-band states, the number of electrons present because of doping is negligible.
45 For a bandwidth-limited pulse, the extrema of the electric field and the zeros of the vector potential always coincide but this is not the case for a pulse with chirp like in our experiment.