Two-dimensional drift–diffusion analysis of magnetic field enhanced THz emission from semiconductor surfaces

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Abstract

A novel two-dimensional drift–diffusion model of carrier dynamics in femtosecond-irradiated n-GaAs and InAs surfaces under a magnetic field is presented. The analysis provides a comprehensive picture of THz emission enhancement mechanisms from semiconductor surfaces. The far-field THz emission is evaluated at various optical excitation levels, magnetic field strengths, and magnetic field orientations. The model accurately describes the power dependence of THz emission from n-GaAs and InAs surfaces for magnetic field values up to ±10 and ±6 T, respectively. Observed saturation effects in both semiconductors for optical excitation densities from 40 nJ/cm² to 2.2 μJ/cm² are in excellent agreement with previously reported experimental data.

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1. Introduction

The generation of ultrashort high power free space THz radiation pulses has become increasingly important for spectroscopic [1] and imaging applications [2,3]. Conventional THz sources include photoconductive (PC) emitters and semiconductor surface emitters. In the former, photoelectrons are accelerated parallel to the surface by biased electrodes, inducing a planar time-varying Hertzian dipole. Although PC emitters provide the highest power, they are difficult to integrate into large area emitters, and the THz power dependence on the biasing field is limited by the dielectric breakdown of the semiconductor (∼10⁵ V/cm). Alternatively, the optical excitation of bare semiconductor surfaces provides a compact source of THz radiation that circumvents electric field breakdown damage, does not require microfabrication, and can be easily integrated into

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large area emitters. In contrast to the planar dipole in PC emitters, photoelectrons in surface emitters are accelerated perpendicular to the surface by the depletion field, thus creating a perpendicular time-dependant dipole. Although surface field values can exceed $10^5$ V/cm, the emission power from surface emitters is typically much smaller than from PC emitters. This is due to polarization dependent refraction at the semiconductor interface, which lowers the radiation efficiency of the perpendicular configuration relative to the planar arrangement. However, by using a magnetic field to tilt the radiating dipole this limitation can be overcome.

Experimental studies have demonstrated a wide range of THz emission enhancement phenomena. Investigations by Corchia et al. [4] into THz emission from femtosecond irradiated n-type GaAs surfaces have shown that emission strength increases with magnetic field to a plateau at 6 T. In undoped InAs samples, McLaughlin et al. [5] have reported a quadratic dependence of THz emission on the magnetic field up to 3 T. Beyond 3 T Ohtake et al. [6] have demonstrated that the radiated THz power saturates at 4 T, decreases to a minimum at 6 T, and increases again at 14 T. Experimental works have also studied enhancement in InAs for various excitation regimes and doping levels. Heyman et al. [7] have reported a lower emission power from n-doped InAs as opposed to undoped InAs, but no explanation of the observed phenomenon was reported. When a magnetic field of 3 T was applied, the emission power from doped InAs was approximately half of that from the undoped sample. In studies by Takahashi et al. [8], the THz emission enhancement factor from undoped InAs was observed to decrease rapidly as the excitation fluence is increased above 0.06 $\mu$J/cm$^2$. For excitations fluences exceeding 0.2 $\mu$J/cm$^2$, the THz power enhancement saturates completely.

Understanding enhancement mechanisms and their dependence on excitation fluence, doping, and magnetic field is crucial to the development of highly efficient THz emitters. To describe magnetic field enhanced emission from InAs and n-GaAs surfaces several models, based on a semiclassical prescription, have been used. Shan et al. [9] attributed enhancement to increased out-coupling of THz radiation from the air–semiconductor interface. Their single particle model, however, does not provide insight into complicated photo-Dember mechanisms and space charge screening effects that are important in high mobility semiconductors such as InAs. To explain THz emission from InAs surfaces, Heyman et al. [7] described the THz emission in terms of classical dynamics of both the electrons accumulated at the surface and in the bulk of the semiconductor. Their picture of electron motion in InAs provides a physical understanding of the enhancement mechanism for a narrow range of magnetic fields below 1.0 T; however, their model breaks down for magnetic field values $>1.0$ T. More recently, Johnston et al. [10] have developed a 3D Monte-Carlo model of THz emission enhancement based on semiclassical electron motion. The model describes emission from GaAs, but for InAs their treatment does not explain the anomalous decrease in THz emission strength near 3 T and fails to predict the dependence of emission on doping concentrations. Moreover, the simulations are computationally intensive and require the accurate input of a large number of parameters related to band structure and scattering rates that are not well known at high magnetic fields.

In this work, we present a novel two-dimensional drift–diffusion model to describe THz emission from n-GaAs and InAs surfaces in a magnetic field. Carrier transport is described in two dimensions to accommodate acceleration induced by the magnetic field. Carrier transport is described in two dimensions to accommodate acceleration induced by the magnetic field. In contrast to single particle semiclassical descriptions, the study provides a novel, consolidated picture of collective carrier behavior during optical excitation, and includes the effects of extrinsic and photogenerated carriers, charge screening effects, and photo-Dember mechanisms. The effects of optical excitation intensity, magnetic field orientation, and magnetic field strength on the far field THz emission are accurately reproduced and in accordance with previously reported experimental observations. The analysis presented provides the most comprehensive model of enhancement to date. The understanding of enhancement phenomena afforded by
our model has potential application in the design of efficient semiconductor THz surface emitters.

2. The model

The geometrical layout of the surface THz emitter used in the analysis is depicted in Fig. 1. An optical pulse is incident on the semiconductor surface at $\theta = 45^\circ$ and the THz emission is detected at $\varphi = 45^\circ$. The magnetic field is oriented along the $y$-axis. Such a configuration has been shown to produce the optimum THz power [11].

In order to accurately describe the carrier dynamics, one must consider the fields existing at the surface of n-GaAs and InAs prior to photoexcitation. Electron trapping at boundary states near the n-GaAs surface causes a surface depletion layer whose electric field, $E_{\text{surf}}(z)$, acts on photoexcited carriers. Here, a Schottky model is used to describe such unperturbed surface field. In this formalism, the electric field is taken to be maximum at the surface ($z = 0$) and decreases linearly away from the surface to zero within the depletion length $w$. The electric field at the surface $E_{\text{surf}}(z)$ and the depletion width $w$ are given by

$$E_{\text{surf}}(z) = \left( \frac{2eN_d\Phi_b}{\varepsilon_0 \varepsilon_r} \right)^{1/2} \left( 1 - \frac{z}{w} \right), \quad \text{for} \quad 0 \leq z \leq w$$

and

$$w = \left( \frac{2e\varepsilon_r \varepsilon_0 \Phi_b}{eN_d} \right)^{1/2},$$

where $N_d$ is the doping density, $\varepsilon_r$ is the relative permittivity of the semiconductor, $\varepsilon_0$ is the permittivity of free space, $e$ is the electron charge, and $\Phi_b$ is the band bending at the surface. In contrast, the origin of the surface field in InAs is due to a different mechanism. In InAs, donor-like surface states and Fermi pinning to the conduction band at the surface results in the accumulation of electrons near the surface [12]. As there is no depletion field to drive the photoexcited electrons, diffusion processes govern the dynamics of photoexcited electrons. The spatial distribution of the steady-state carrier density, $n_{\text{ss}}$, beneath the interface is estimated from experimental data using $n_{\text{ss}} = n_0 e^{-\beta z}$, where $\beta = 30$ nm and $n_0$ is the equilibrium electron concentration [13,14].

In semiconductors at high carrier densities ($>10^{15}$ cm$^{-3}$) and in the time scales of interest ($\sim 100$ fs) where carrier–carrier, carrier–phonon, and carrier–impurity scattering rates are high, the Boltzmann transport picture describes photoexcited carrier motion accurately. In this formalism, the holes are taken to be stationary and electron acceleration is governed by an electrostatic potential, $\psi(x,z,t)$. This electrostatic potential incorporates the potential due to the surface field and the potential due to space charge separation. The current density, $\vec{j}(x,z,t)$, is obtained from the momentum conservation equations under relaxation

$$\frac{\partial \vec{j}(x,z,t)}{\partial t} = \frac{e}{m^*} \left( en(x,z,t) \vec{\nabla} \psi(x,z,t) + \vec{j}(x,z,t) \times B \right)$$

$$+ \frac{eD}{\mu} \frac{\partial n(x,z,t)}{\partial t},$$

where $B$ is the static magnetic field, $m^*$ is the electron effective mass, $D$ is the electron diffusion coefficient, $\mu$ is the electron mobility, and $\tau$ is the steady-state momentum relaxation time due to carrier collisions.

To satisfy conservation of charge flux, the current density is related to the carrier density through the continuity equations

$$\frac{\partial n(x,z,t)}{\partial t} = \frac{\varepsilon(1-R)}{\hbar \omega} I(x,z,t) + \frac{1}{e} \vec{\nabla} \cdot \vec{j}(x,z,t)$$

$$- \frac{n(x,z,t)}{\gamma}$$

Fig. 1. Configuration of the laser pulse and THz emission with respect to the semiconductor surface.
and

$$\frac{\partial p(x,z,t)}{\partial t} = \frac{z(1-R)}{h\omega} I(x,z,t) - \frac{p(x,z,t)}{\gamma},$$  \hspace{1cm} (5)

where $p(x,z,t)$ is the hole density, $h$ is Planck’s constant, and $\omega$ is the frequency of the laser pulse, $R$ is the reflectivity of the semiconductor, $\gamma^{-1}$ is the carrier lifetime, and $z^{-1}$ is the semiconductor absorption depth. The time dependent intensity, $I(x,z,t)$, of the excitation pulse is given by

$$I(x,z,t) = I(t) \exp\left(-\frac{(x \cos \theta - z \sin \theta)^2}{\sigma^2} - (x \sin \theta + z \cos \theta)\tau\right),$$ \hspace{1cm} (6)

where $I(t)$ is the Gaussian temporal envelope of the pulse intensity, $\theta$ is the angle of incidence of the laser pulse with respect to the surface, and $\sigma$ is the effective pulse spot diameter. Through the application of Poisson’s equation,

$$\nabla^2 \psi(x,z,t) = \frac{e}{\varepsilon_2 \varepsilon_0} \left(n(x,z,t) - p(x,z,t) + N_d\right)$$  \hspace{1cm} (7)

charge separation effects induced by carrier transport are coupled self-consistently to the electrostatic potential.

The system of coupled partial differential Eqs. (3), (4), (5) and (7) are solved for all time using a predictor corrector method where at each step, $\psi(x,z,t)$ is determined. An over-relaxed Gauss–Siedel method is utilized to evaluate $\psi(x,z,t)$ and solve self-consistently for $j(x,z,t)$. At $z = 0$, Neumann boundary conditions are imposed on $\psi(x,z,t)$ and at $|x| > |\sigma|$ and $z > z^{-1}$, $\psi(x,z,t)$ is set to zero. The emitted THz electric field components originating from the focal volume is given by

$$E_{x,z}(t) = \frac{t_{x,z}(\varphi)\sigma_y}{4\pi\varepsilon_0 \varepsilon_r c^2} \frac{\partial}{\partial t} \int \frac{f_{x,z}(x',z',t)}{|L - r'|} \, dV,$$ \hspace{1cm} (8)

where $L$ is the distance of the detector from the origin, $c$ is the speed of light, $\sigma_y$ is the pulse spot size in the $y$-direction, and $t_{x,z}(\varphi)$ incorporates the transmission coefficient due to refraction at the air–semiconductor interface. The transmission coefficients for the THz field components parallel $E_x(t)$ and perpendicular $E_z(t)$ to the surface are expressed as \cite{9}

\begin{equation}
\hspace{1cm}
t_x(\varphi) = \frac{(\varphi - \sin^2 \varphi)^{1/2}}{\varphi \cos \varphi + (\varphi - \sin^2 \varphi)^{1/2}},
\end{equation}

\begin{equation}
\hspace{1cm}
t_z(\varphi) = \frac{\sin \varphi}{\varphi \cos \varphi + (\varphi - \sin^2 \varphi)^{1/2}},
\end{equation}

respectively.

3. Results and discussion

The two-dimensional drift–diffusion model is then used to investigate the carrier dynamics in n-doped GaAs and undoped InAs for varying magnetic field strengths and excitation regimes. The results are compared with previously reported experimental data. The parameters used in the simulation are listed in Table 1.

To explain the complex carrier dynamics, both the current density and electrostatic potential must be considered simultaneously. Fig. 2 depicts the sequence of events of the spatio-temporal evolution of the current density, $j(x,z,t)$, in n-GaAs for $B = 0$ T at times $t = -100, 300,$ and $700$ fs after photo excitation with a $100$ fs, $800$ nm laser pulse. The corresponding electric potential is plotted in Fig. 3. Prior to the arrival of the pulse at $t = -100$ fs, the steady-state carrier dynamics govern the potential distribution. The equilibrated steady-state carrier population sets up a depletion field extending 500 nm into the n-GaAs surface with a peak value of 7.6 kV/cm. At $t = 0$ fs a high density of carriers within the depletion region has been photoexcited by the laser pulse. At

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<th>Table 1 Parameters used for InAs and n-GaAs</th>
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$t = 300 \text{ fs}$, the depletion field within 100 nm of the surface has been almost entirely screened by the photogenerated carriers as evidenced by the zero gradient of the electrostatic potential in Fig. 3. Coupled drift and diffusion forces govern the dynamics of photogenerated carriers. As a result, electron diffusion from the high carrier concentration region causes a small current density component parallel to the surface, $j_x(x, z, t)$. Meanwhile, the current density component perpendicular to the surface $j_z(x, z, t)$ is driven by the strong surface depletion field and exceeds $j_x(x, z, t)$ by a factor of 12, as shown in Fig. 2. Evidently, diffusion currents play a minor role in the current density evolution in n-GaAs. This prediction is consistent with observations by Johnston et al. [10]. At longer times, $t = 700 \text{ fs}$, carrier scattering and carrier recombination reduce current densities to equilibrium values.

When the magnetic field is non-zero, the carrier dynamics are significantly modified. Fig. 4 illustrates the evolution of the current density in n-GaAs when a magnetic field of $B = 6 \text{ T}$ is applied. In contrast to the $B = 0 \text{ T}$ case, when a magnetic field is applied $j_y(x, z, t)$ does not exhibit a bipolar distribution at $t = 300 \text{ fs}$. Instead, carriers that are initially driven perpendicular to the surface by the depletion field are accelerated into the parallel direction by the magnetic field, enhancing the current density component parallel to the surface. In this situation, the current density perpendicular to the surface, $j_z(x, z, t)$, is only a factor of 2 larger than $j_x(x, z, t)$. As $j_y(x, z, t)$ is enhanced, the coupling efficiency of the THz emission increases due to

![Fig. 2. Spatio-temporal evolution of the perpendicular (left) and parallel (right) current densities within n-GaAs at different times for $B = 0 \text{ T}$](image1)

![Fig. 3. Electrostatic potential as a function of depth in n-GaAs at various times for $B = 0 \text{ T}$](image2)

![Fig. 4. Spatio-temporal evolution of the perpendicular (left) and parallel (right) current densities within n-GaAs at different times for $B = 6 \text{ T}$](image3)
polarization sensitive refraction at the semiconductor–air interface.

On the other hand, carrier dynamics are drastically different in InAs due to the higher electron diffusivity and the absence of a surface depletion field. Again, to illustrate the carrier dynamics, both the current density and electrostatic potential are considered. Fig. 5 depicts the spatial dependence of the electrostatic potential beneath the InAs surface at various times for $B = 0$ T. Prior to excitation at $t = -100$ fs, the constant potential within the semiconductor describes the absence of a surface field. Upon the photoexcitation of electron–hole pairs at $t = 0$ fs, diffusion of high mobility electrons will cause charge separation and the establishment of a built-in photo-Dember field (2 kV/cm). This can be illustrated later at $t = 300$ fs, where a large photo-Dember potential is established within 50 nm of the surface. At $t = 700$ fs, the photo-Dember field penetrates a maximum depth of $\sim 100$ nm into the InAs surface. The photo-Dember field vanishes for times greater than the carrier lifetime. While the photo-Dember field is established, ultrafast electron diffusion drives large currents within the semiconductor. Fig. 6 depicts the current density components in undoped InAs at times $t = -100$, 300, and 700 fs for a zero magnetic field strength. At $t = -100$ fs, few carriers are excited by the weak wing of the Gaussian pulse, resulting in insignificant, but finite, current densities in both the $x$- and $z$-directions. At the peak of the pulse, $t = 0$ fs, large carrier populations have been created by the absorption of the laser pulse within 143 nm of the InAs surface. As shown for $t = 300$ fs in Fig. 6, carrier diffusion, which is proportional to the carrier density gradient in the photoexcited region, results in the unipolar and bipolar spatial distributions of $j_x(x,z,t)$ and $j_z(x,z,t)$, respectively. At a later time at $t = 700$ fs, carrier scattering and recombination drive the current densities to equilibrium values.

When a high magnetic field is applied, the Lorentz force alters the carrier trajectories and thus the spatial distribution of the current densities. This is illustrated in Fig. 7 by following the current density spatio-temporal evolution for $B = 3$ T. In contrast to the $B = 0$ T case, at $t = 300$ fs, the application of the magnetic field has altered the current density distributions such that $j_x(x,z,t)$ is now bipolar, whereas $j_z(x,z,t)$ is unipolar. This is equivalent to the reorientation of the electrical dipole by $90^\circ$. When compared to n-GaAs, the magnitude of $j_x(x,z,t)$ relative to $j_z(x,z,t)$ in InAs has not been significantly enhanced with the application of a magnetic field. Instead, the experimentally observed THz radiation emission enhancement in InAs can be attributed to the modified current density spatial distributions. That is, for $B = 0$ T, the bi-directional nature of $j_x(x,z,t)$ results in far field cancellation of the THz emission.
component parallel to the surface. However, when \( B = 3 \) T, \( j_x(x,z,t) \) is uni-directional, resulting in the coherent addition of the emitted THz radiation in the far-field. This suggests that the observed THz radiation enhancement in InAs is primarily ascribed to the spatial reorientation of the current densities.

The THz radiation emission from the surface of n-GaAs is shown in Fig. 8 for various magnetic field strengths. For \( B = 0 \) T, the THz pulse component polarized along the x-axis is lower by a factor of \(~250\) and is relatively broader than the z-polarized field component. This is due to the fact that \( j_x(x,z,t) \) is driven by slow diffusion processes while \( j_z(x,z,t) \) originates from rapid carrier acceleration by the surface field. When the applied magnetic field is increased, the amplitude of \( E_x(t) \) increases significantly while \( E_z(t) \) slightly decreases. This results in an overall enhancement of the THz emission amplitude. In addition, it is observed that the electric field components, \( E_x(t) \) and \( E_z(t) \), evolve from a predominantly unipolar pulse envelope at low magnetic field into a pronounced bipolar pulse shape as the field increases. In the spectral domain, this behavior manifests itself through spectral broadening and shifting. The emission spectra at \( B = 3 \) and 6 T are depicted in Fig. 9. At \( B = 3 \) T, the spectrum consists of a single broad peak centred at 0.6 THz. As the magnetic field reaches 6 T, the spectrum broadens and the spectral peak shifts towards 1 THz. This behavior is in excellent agreement with experimental results demonstrated by Corchia et al. [1] for n-GaAs using similar experimental parameters. In their experiment, similar spectral broadening and shifting is observed as the magnetic field is increased from 2 to 4 T.
The temporal and spectral behavior of the THz emission from InAs are also analyzed as a function of the applied magnetic field. The components $E_x(t)$ and $E_z(t)$ from undoped InAs are presented in Fig. 10 for varying magnetic fields. At $B = 0$ T, both $x$- and $z$-polarized components have broad, bipolar pulse shapes due to diffusion-driven current densities. As the magnetic field is increased to 3 T, the THz pulse exhibits temporal narrowing and significant post-peak oscillations following photoexcitation. Further increase in the magnetic field up to 12 T results in overdamped emission with reduced amplitude. Interestingly, these temporal characteristics are similar to what was experimentally observed by McLaughlin et al. [5]. The emission spectra at $B = 3$ and 12 T are illustrated in Fig. 11. At $B = 3$ T, the spectrum is broad and centred at 3.25 THz. When the magnetic field is increased to 12 T, the spectrum shifts toward lower frequencies and peaks at approximately 0.75 THz. Although the locations of the peaks reported in [6] occur at lower frequencies, the overall predicted behavior agrees with the experimental emission spectra. In their work, a marked shift in the THz spectrum toward lower frequencies is observed as the field is increased from 3 to 14 T.

An important prediction from the model is the THz radiation power dependence on the applied magnetic field. Fig. 12 illustrates the power of the THz emission from n-doped GaAs as a function of the magnetic field strength at a fluence of 0.22 $\mu$J/cm$^2$. Interestingly, the magnetic field dependence of the THz emission power is asymmetric about 0 T with a minimum THz power emission occurring at approximately 0.3 T. Experimentally, Heyman

![Fig. 10. The z-polarized (top) and x-polarized (bottom) temporal waveforms of THz emission from InAs at various magnetic field strengths.](image)

![Fig. 11. Normalized THz emission spectra from InAs at $B = 3$ and 12 T.](image)

![Fig. 12. THz emission power for: (a) n-GaAs and (b) InAs as a function of the magnetic field strength.](image)
et al. [7] have measured a similar trend where an offset in the minimum THz power was observed at 0.9 T. The discrepancy between the experiment and the model are likely due to small differences in the optical pulse fluence and the doping concentrations of the sample. From our model, the offset in the minimum is attributed to off-normal illumination of the surface. As there are significant diffusion currents along the $x$-direction at $B = 0$ T, the minimum THz emission does not occur at $B = 0$ T. Instead, when a slight magnetic field of 0.3 T is applied, re-direction of a small amount of $j_z(x,z,t)$ to the $x$-direction results in minimum THz radiation power due to cancellation of the emission along $\pm x$. Beyond a magnetic field magnitude of $|1|$ T, the THz emission increases linearly and then saturates at $\sim |5.5|$ T. Such a behavior is also in accordance with experimental studies reported by Heyman et al. [7].

The magnetic field dependence of the THz emission power from n-InAs at a fluence of 0.04 $\mu$J/cm$^2$ is illustrated in Fig. 12. The emission power increases rapidly from 0 T and peaks at approximately 3 T. The peak power at +3 T exceeds that at $-3$ T, consistent with several experimental studies [6,7,15]. The asymmetric power dependence on the magnetic field can again be explained from our model by considering the effect of off-normal photo-excitation. In such a configuration, the diffusion forces adjacent to the photoexcited region are not symmetric. As a result, the amount of dipole rotation induced by the magnetic field will depend on the direction of rotation and consequently, the direction of the magnetic field. Unlike previous Monte-Carlo simulations, our analysis accurately models the decrease in the emission power beyond 3 T in InAs. However, the model is unable to reproduce the recovery in the THz emission reported above $B = 6$ T [6].

The optical excitation fluence is a crucial factor in the magnetic field-induced enhancement of THz emission. Fig. 13 illustrates the THz emission power versus excitation fluence up to 2.2 $\mu$J/cm$^2$. For fluences from 40 nJ/cm$^2$ to 0.5 $\mu$J/cm$^2$, the power enhancement is approximately 25. This enhancement value predicted in the model agrees well the reported experimental enhancement factors (10–35) using n-GaAs [7,16]. This behavior is in accordance with studies by Corchia et al. [1], where the enhancement is nearly constant for fluences from 10 to 200 nJ/cm$^2$. Interestingly, as the fluence is increased above 0.5 $\mu$J/cm$^2$, carrier screening of the surface field results in a decrease in the emission enhancement factor. The screening, which is directly proportional to the energy fluence, causes a linear decrease in the THz emission power as the energy fluence increases. Fig. 13(b) depicts the enhancement factor in InAs versus optical excitation fluence. For a fluence of 0.25 $\mu$J/cm$^2$, an enhancement factor of 84 is predicted. This is in good agreement with the enhancement factor of 74 measured in experiments by Heyman et al. [7], where 60 nJ/cm$^2$ pulses and an identical geometry are used. As the excitation fluence increases above $0.5$ $\mu$J/cm$^2$ the maximum power enhancement decreases markedly, saturating above $0.9$ $\mu$J/cm$^2$. The decreased enhancement with increasing fluence is attributed to charge screening of the photo-Dember field, limiting the dipole rotation induced by the magnetic field. A much sharper decrease is observed in InAs as opposed to n-GaAs since diffusion...
and photo-Dember effects more responsive to the carrier densities than drift effects. The predicted fluence dependence is found to be in excellent agreement with experimental data by Takahashi et al. [8] plotted as an inset of Fig. 13(b). To study the effects of doping on THz emission from InAs, doped InAs is modeled by simply increasing the background concentration of electrons in InAs while keeping all other parameters constant. For the doped sample \( n_o = 1 \times 10^{16} \text{ cm}^{-3} \), the maximum THz emission power is predicted to decrease by a factor of 1.4. This value is in good agreement with [7] where the emission from the doped InAs sample \( n_o = 7.5 \times 10^{16} \text{ cm}^{-3} \) was 2 times weaker than the nominally undoped InAs \( n_o = 1.8 \times 10^{16} \text{ cm}^{-3} \). The decreased emission power from the doped sample is attributed to the increased charge screening due to the increased steady-state electron accumulation layer.

4. Conclusion

In conclusion, we have investigated THz emission for n-GaAs and InAs surfaces based on drift–diffusion transport model. Unlike previous semi-classical and Monte-Carlo simulations, the current analysis provides a complete two-dimensional picture of current densities and carrier population evolution and includes charge screening, photo-Dember, and THz refraction effects. The work clearly demonstrates that diffusion and photo-Dember driven current densities are responsible for THz emission from InAs, whereas drift currents dominate in n-GaAs. THz power enhancement and spectra for n-GaAs are well described for magnetic fields up to 10 T. Anomalous decrease in the THz emission power from InAs surfaces above 3 T is accurately reproduced. In addition, experimentally observed asymmetric power dependences on the magnetic field in InAs and n-GaAs have been explained and in both cases, is ascribed to off-normal surface illumination. For fluences up to 2.2 \( \mu \text{J/cm}^2 \), the enhancement as a function of fluence for n-GaAs and InAs is in excellent agreement with previously reported experimental data. The analysis provides a physical picture and a novel theoretical understanding of enhancement of THz emission from semiconductor surfaces. The application of this model will certainly be used in the design and optimization of THz surface emitters.

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