Origin of magnetic field enhancement in the generation of terahertz radiation from semiconductor surfaces

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We present a theory of the magnetic field enhancement of terahertz (THz) emission from photogenerated carriers in the surface depletion region of a semiconductor. A combination of the Drude–Lorentz model for the carrier dynamics with an appropriate solution of the radiation problem is sufficient to explain the strong B-field enhancement in THz radiation that has been observed experimentally. The effect arises primarily from the increased radiation efficiency of transient currents flowing in the plane of the surface. The model provides quantitative agreement with experiment for the pronounced angular dependence of the enhancement and predicts the correct trend for the enhancement in a variety of materials. © 2001 Optical Society of America

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The motion of charge carriers produced in the depletion region of a semiconductor surface by an ultrashort laser pulse provides a convenient source of coherent electromagnetic radiation in the terahertz (THz) spectral region. It was demonstrated as early as 1993 (Refs. 1 and 2) that the application of an external magnetic field could significantly enhance this emission. Subsequent investigations have characterized this effect extensively, including its variation as a function of the strength and orientation of the B field,[1–9] of the excitation energy and geometry,[10,11] of the photoconductive material,[5] and of the temperature.[6–9]

Surprisingly, however, no complete model has been developed that is capable of predicting the large enhancements that have been seen experimentally, although the role of the magnetic component of the Lorentz force in modifying the photocurrent was identified in the initial report and most subsequent works. The natural starting point for such an explanation is a semiclassical model of carrier dynamics. Within this picture, one can, however, identify what appears to be a fundamental difficulty: The magnetic force does not increase the magnitude of the acceleration of a charge moving in a uniform electric field. Thus, one may reason, the radiated THz field, which is proportional to the carrier acceleration, cannot increase.

In this Letter we present an analysis of the THz emission process within the Drude–Lorentz model. We show, in contrast with the argument given above, that the THz emission can indeed be significantly enhanced by an external B field. The critical factor is the radiation efficiency for current transients flowing in differing directions. Because of the strong dielectric screening that is present in photoconductive emitters, current flowing parallel to the surface radiates much more efficiently than an equivalent current moving perpendicular to the surface. The enhancement consequently arises when the Lorentz force produces currents in the plane of the surface that would otherwise flow along the surface normal. Thus, the enhanced emission does not require an increase in the *magnitude* of the carrier acceleration but only a change in its *direction*. To test the theory we have obtained experimental data on B-field enhancement as a function of the emission angle of the THz radiation. Good agreement is found. In addition, within the model, the B-field enhancement is determined by the carrier mobility and the optical properties of the photoconductive medium. These simple predictions are shown to reproduce trends that were observed experimentally for several distinct material systems and to provide semiquantitative agreement for the magnitude of the enhancement. The analysis also predicts the observed dependence of the enhancement on the B-field strength for weak fields and the recently reported saturation behavior.[4]

We first present the model. Consider THz radiation generated by illumination of a semiconductor surface with an ultrashort laser pulse that is incident at an angle θ. If the spot size of the laser is appreciably larger than a typical THz wavelength, the emission will be peaked along the reflected direction. We find for the $\hat{s} = \hat{y}$ and $\hat{p} = -\cos\theta\hat{x} + \sin\theta\hat{z}$ polarized components of the electric field in the far field

$$E_s(t) = -\frac{2V}{c^2 r}\frac{1}{\cos\theta + (\epsilon - \sin^2\theta)^{1/2}}\left(\frac{\partial \mathbf{j}}{\partial t}\cdot\hat{y}\right),$$

$$E_p(t) = -\frac{2V}{c^2 r}\frac{\sin\theta}{\epsilon\cos\theta + (\epsilon - \sin^2\theta)^{1/2}} \times\left(\frac{\partial \mathbf{j}}{\partial t}\cdot\hat{z} - \gamma\frac{\partial \mathbf{j}}{\partial t}\cdot\hat{x}\right). \quad (1)$$

Here V denotes the volume in which the current density, $\mathbf{j}(t)$, is present; r is the distance of propagation

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for the center of the THz beam; and ϵ is the THz dielectric constant of the sample. $\partial \mathbf{j}(t)/\partial t$ is evaluated at retarded time $t' = t - r/c$. The parameter $\gamma = (\epsilon - \sin^2\theta)^{1/2}/\sin\theta$, which relates to the radiation efficiency, will be discussed below.

In addition to this electromagnetic description of the radiation process, we need a model for the current density under pulsed photoexcitation. To this end, we make use of the following relation for the current induced by a time-varying carrier density, $N(t)$, in the static depletion field $\mathbf{E}$ (Ref. 12):

$$\mathbf{j}(t) = \int_{-\infty}^{t} eN(t')\overleftrightarrow{\mu}(t - t')\mathbf{E}\mathrm{d}t'. \qquad (2)$$

Here the time-domain mobility tensor, $\overleftrightarrow{\mu}(t)$, describes the linear response of the carriers for the static electric field but includes the influence of the applied B field parametrically to all orders. Physically, $\mu_{ij}(t)$ represents the ith component of the acceleration $\mathbf{a}(t)$ of a charge carrier that started at rest at $t = 0$ under the influence of a unit electric field along the jth direction.

Equations (1) and (2) predict the THz emission in terms of the material response, $\overleftrightarrow{\mu}(t)$, and the carrier density, $N(t)$. To simplify the analysis we suppose that the laser-pulse duration $<<$ carrier scattering time $\ll$ carrier lifetime. We may then approximate $N(t)$ as a step function, $\theta(t)$, and rewrite Eq. (2) as

$$\partial \mathbf{j}(t)/\partial t = eN_0\theta(t)\mathbf{a}(t). \qquad (3)$$

Here $\mathbf{a}(t) = \overleftrightarrow{\mu}(t)\mathbf{E}$ and N_0 is the carrier density created by the laser excitation pulse. Equations (1) then imply that the THz field is proportional to the carrier acceleration.

The motion of the charge carriers is treated within a Drude–Lorentz approximation, in which the carrier velocity is governed by $\mathrm{d}\boldsymbol{\nu}/\mathrm{d}t + \boldsymbol{\nu}/\tau = (e/m^*)[\mathbf{E} + (\boldsymbol{\nu}/c) \times \mathbf{B}]$. Here m^* and τ are, respectively, the effective mass and the mean scattering time of the charge carriers. The depletion field is denoted $\mathbf{E} = E\hat{z}$, and the external magnetic field $\mathbf{B}$. Since the component of $\mathbf{B} \mathbin{//} \mathbf{E}$ has no effect on the motion of the carriers, we take $\mathbf{B} = B\hat{y}$ (perpendicular to the plane of incidence) as a representative configuration with $\mathbf{B} \perp \mathbf{E}$. Imposing the initial condition $\boldsymbol{\nu}(t = 0) = 0$, we then find for the carrier acceleration

$$\mathbf{a}(t) = eE/m^* e^{-t/\tau}(-\sin\omega_c t\hat{x} + \cos\omega_c t\hat{z}), \qquad (4)$$

where $\omega_c = eB/m^*c$ is the cyclotron frequency.

For this configuration, the s-polarized radiation vanishes ($E_s \equiv 0$), and the p-polarized THz field is given by

$$E_p(t) \sim \theta(t)e^{-t/\tau}(\cos\omega_c t + \gamma\sin\omega_c t)\sin\theta. \qquad (5)$$

The significance of the radiation factor γ now becomes clear. As can be seen from Eq. (4), the B field does not change the magnitude of the acceleration, $|\mathbf{a}|$. However, the B field converts some of the current initially flowing along the surface normal ($\hat{z}$) to the component parallel to the surface ($\hat{x}$), which, according to relation (5), radiates γ times more efficiently. Since $\gamma = (\epsilon - \sin^2\theta)^{1/2}/\sin\theta \gg 1$ for typical photoconductive materials (with $\epsilon \gg 1$), THz emission is enhanced by the B field.

Let us explore the enhancement more quantitatively by defining a power-enhancement factor $\eta_P = \int_{-\infty}^{+\infty}\mathrm{d}t|E_B(t) - E_{B=0}(t)|^2/\int_{-\infty}^{+\infty}\mathrm{d}t|E_{B=0}(t)|^2$ for the THz emission. From relation (5) we obtain

$$\eta_P(\gamma, x = \mu_0 B) = \frac{\gamma^2}{2} + \frac{3}{2} + \left(\frac{1}{2} - \frac{\gamma^2}{2} + \gamma x\right) \Big/ (1 + x^2) - (2 + \gamma x)/(1 + x^2/4), \qquad (6)$$

which is determined solely by the radiation factor γ and the product $x = \mu_0 B = \omega_c\tau$ of the B-field strength and the carrier mobility, $\mu_0 = e\tau/m^*$. As an example, we show in Fig. 1 the enhancement for an InAs emitter excited by a pump laser at 45°. A strong enhancement is predicted for the actual THz dielectric constant $\epsilon = 14.6$ ($\gamma^2 = 28$), but little effect is expected under the same circumstances for a material with $\epsilon = 1$.

The asymptotic behavior of the enhancement as a function of the B-field strength emerges readily from Eq. (6): In the regime of weak B fields ($x \ll 1$), the enhancement is quadratic in the B-field strength, with $\eta_P \approx (\gamma\mu_0)^2 B^2/2$. For strong B fields the enhancement saturates at $\eta_P = \gamma^2/2 + 3/2$, a value that depends only on the optical properties of the material, not on the carrier mobility. The onset of saturation occurs for $x = x_s \approx 6/\gamma$ or $B_s \approx 6/\gamma\mu_0$. Thus materials with higher carrier mobilities will reach saturation for weaker B fields.

To test this model experimentally, we chose to probe the angular dependence of the enhancement. This quantity has a strong and distinctive variation, which is largely independent of the detailed model of the material response. The experimental apparatus was similar to that used in earlier experiments.[5] With the configuration of the pump laser and the magnetic field as described above, we measured the waveforms of THz radiation, $E(t)$, without and with the external B field to determined the power-enhancement factor,

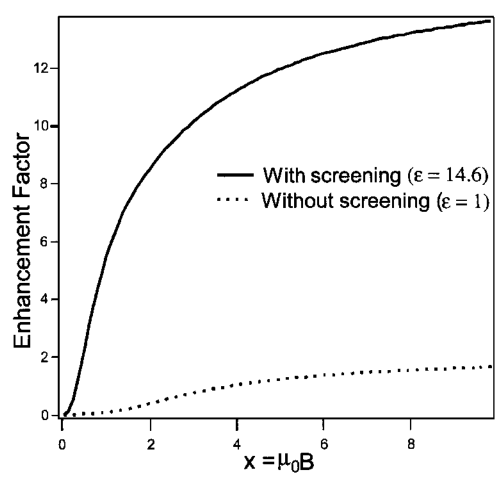

Fig. 1. Magnetic field dependence of the THz enhancement predicted for an InAs emitter excited by a pump laser at 45° (solid curve) with and (dotted curve) without the inclusion of the THz dielectric response.

Table 1. Comparison of B-Field Enhanced THz Emission in Five Lightly Doped Semiconductors

Property	Semiconductor				
	GaSb	InP	GaAs	InAs	InSb
THz dieletric constant ϵ^a	15.7	12.4	13.2	14.6	17.7
Bulk electron mobility μ_b at 300 K ($cm^2/V\ s$)[a]	5000	5000	9000	33,000	78,000
THz power enhancement for weak B fields (T^{-2})	0.7	3.7	11.1	8.8	62.0
Inferred carrier mobility μ_0 ($cm^2/V\ s$)[b]	2200	5600	9300	8000	19,000

[a]Parameters from Ref. 13.
[b]The effective carrier mobility μ_0 is determined so that the model matches the experimental THz power-enhancement factors reported in Ref. 5.

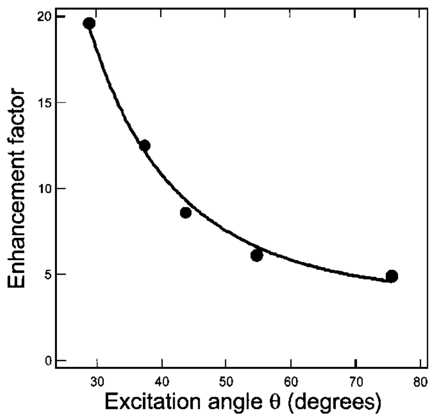

Fig. 2. Experimental data for the angular dependence of the THz enhancement of an InAs sample under a magnetic field of 1 T (filled circles). The solid curve corresponds to a fit to the model presented in the text.

η_P. We illustrate in Fig. 2 the angular dependence of the (unsaturated) enhancement for an InAs emitter under a magnetic field $B = 1$ T. The enhancement can be seen to vary sharply with the excitation angle, θ. The experimental data agree very well with a fit of the model. Note that, in the limit of $\epsilon \gg 1$, the theory predicts a universal angular dependence of $\eta_P \propto \gamma^2 = |\epsilon - \sin^2\theta|/\sin^2\theta \sim 1/\sin^2\theta$.

We also applied the model to analyze the previously reported enhancement in a series of III–V semiconductors.[5] The quadratic dependence of η_P predicted above was seen in all samples at low B-field strengths. As for the magnitude of the enhancement, Table 1 summarizes the relevant material properties (THz dielectric constant and bulk electron mobility μ_b at 300 K) and the measured enhancement coefficient. We inferred the carrier mobility, μ_0, required within our model to match the experimental enhancement. Comparing μ_0 and μ_b, we conclude that the model predicts the correct trend for all the materials studied. The discrepancy between μ_0 and μ_b generally corresponds to an effective mobility (μ_0) that is less than the bulk carrier mobility (μ_b). This situation is most pronounced for the materials with high mobilities (InAs and InSb). It may reflect additional scattering in the surface region and the inability of the carriers to reach their full steady-state velocity within the depletion region of the material, effects that might be manifested in other surface transport properties.

We have developed a simple but complete semiclassical treatment of the magnetic field enhancement of THz emission from semiconductors. Combining the Drude–Lorentz model with radiation theory, we were able to predict the asymptotic behavior of the phenomenon with B-field strength and the correct trend for the magnitude of the enhancement in various material systems. The differing radiation efficiency for transient currents flowing parallel and perpendicular to the surface was identified as the principal source of the enhancement. The importance of this effect was demonstrated in experimental measurements of the angular dependence of the enhancement factor.

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