Breakdown of Magnetic Insulation in Semiconductor Plasmas

K. Papadopoulos, A. Zigler, D. L. Book, C. Cohen, and D. Hashimshony

Abstract— **A theoretical analysis of the current response of strongly magnetized electrically biased photoconductors to short laser pulses, with emphasis on the breakdown of magnetic insulation, is presented. There are two regimes that result in breaking of the magnetic insulation during the "on" time of the pulse: 1) the collisionless regime, applicable to pulses with duration $\tau_0 < 1/\nu$, where ν is the collision frequency, in which the magnetic insulation is broken by a polarization-like current induced by the fast rate of increase of the carriers, and 2) the collisional regime, applicable to pulses with $\tau_0 > 1/\nu$, where the magnetic insulation is broken at high carrier density due to the nonlinear dependence of the collision frequency on the carrier density. A simple experiment was performed which confirms the physics of the collisional regime. It is shown that the presence of the magnetic field can significantly reduce the response time of photoconductors. Response times shorter than a picosecond can be achieved in the collisionless regime.**

I. Introduction

The interaction of short optical pulses with electrically biased photoconductors is a frontier scientific topic with important technological implications [1]–[3]. For example, the temporal response of a semiconductor plasma to short laser pulses in the presence of electric bias controls the development of subpicosecond optoelectronic devices, switches, and microwave pulses with terahertz bandwidths [4]–[6]. In these optoelectronic devices, the transition from the conductive to the insulating state is controlled purely by the time scales for creation and loss of high-mobility carriers in the semiconductor plasma. While the creation of carriers by a short laser pulse is an intrinsically fast process, their loss, which is controlled by either recombination or carrier sweep-out, is a slow process resulting in relatively long and undesirable switch-off transients. In this paper we present a theoretical and experimental study of the response of a magnetized semiconductor plasma to an optical pulse of length τ_0. New physical effects associated with the strong magnetization condition $\Omega/\nu \equiv \mu B > 1$, where Ω and ν are the carrier cyclotron and collision frequencies, μ is the mobility, and B is the magnetic field strength, result in transitions between the insulating and conducting states on time scales much faster than for the unmagnetized case. This paper concludes with a brief discussion of the applications of the physical model to photoconducting semiconductor switches (PCSS) and to subpicosecond photonic technology.

The novel physical aspects of the magnetized response can be illustrated by considering a simple model that describes the electron response of an electrically biased semiconductor to a laser pulse that generates electron–hole pairs (exciting pulse). Incorporation of the hole response is straightforward. Each charge passes through the same phases, but the ones created later are at an earlier stage in their evolution. The temporal response of the electron current $\mathbf{J}(t)$ to an exciting laser pulse in a photoconductor is therefore an integral over successive increases $dn = \dot{n}dt$ in the population. In the presence of an electric field $\mathbf{E} = \mathbf{e}_x E$ and a magnetic field $\mathbf{B} = \mathbf{e}_z B$ its evolution is described by the set of equations

$$\mathbf{J}(t) = e\int_0^1 dt'\dot{n}(t')\boldsymbol{v}(t-t') \tag{1}$$

$$\dot{\boldsymbol{v}}(t) = \frac{e}{m}(\mathbf{E} + \mathbf{v}\times\mathbf{B}) - \nu\boldsymbol{v} \tag{2}$$

$$\dot{n}(t) = Q(t) - n/\tau. \tag{3}$$

Here, $n(t)$ is the electron density, $Q(t)$ is the volume production rate of carriers by the laser pulse, and τ is the carrier loss time. We have assumed for simplicity that the electric field is limited to values small enough that the mobility is independent of the electric field.

The key concept of magnetic insulation and the processes leading to its breakdown in photoconducting media can be illustrated by solving (1)–(3) in the absence of carrier losses, i.e., for times $t \ll \tau$. If we assume that the initial velocity of the carriers is negligible ($\mathbf{v} = 0$), their velocity in the direction of the electric field as a function of time is given by

$$\mathbf{v}(t) = \frac{\mu\mathbf{E}}{1+\mu^2B^2}\{1 + \exp(-\nu t)[\mu B\sin(\Omega t) - \cos(\Omega t)]\} \tag{4}$$

and the associated current by

$$\mathbf{J}(t) = \mathbf{J}_c(t) + \mathbf{J}_p(t) \tag{5a}$$

$$\mathbf{J}_c(t) = \frac{n_0 e\mu\mathbf{E}}{1+\mu^2B^2} \tag{5b}$$

$$\mathbf{J}_p(t) = \mathbf{J}_c(t)\frac{\exp(-\nu t)}{n_0}\int_0^t dt'Q(t')\exp(\nu t') \times \{\mu B\sin[\Omega(t-t')] - \cos[\Omega(t-t')]\}$$

Manuscript received July 19, 1995; revised March 11, 1996. This work was supported by Advanced Power Technologies, Inc. internal funds.

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Publisher Item Identifier S 0093-3813(96)05204-6.

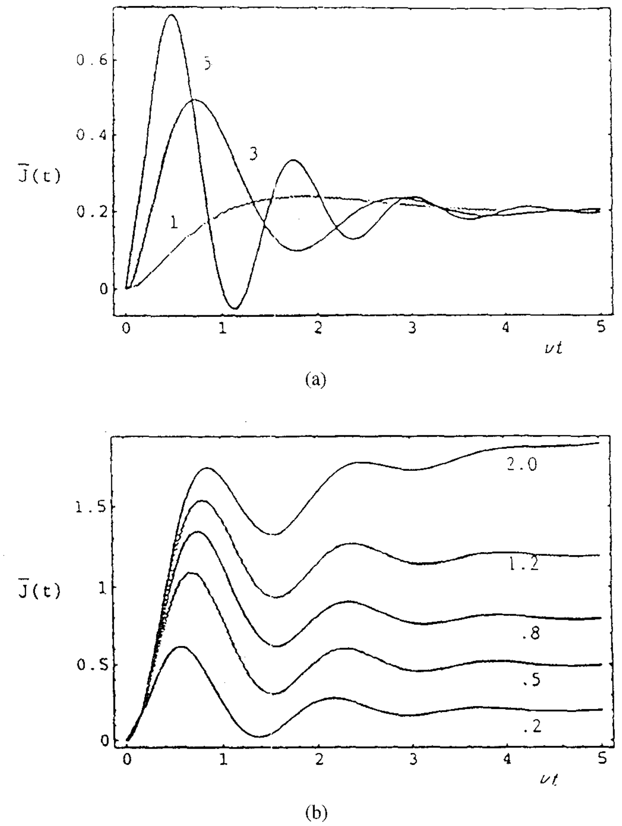


Fig. 1. Temporal response of normalized current density $\bar{J}(t) = \frac{J(t)}{J}$ in the short-pulse case for (a) $\nu\tau_0 = 0.2$ and $\mu B = 1, 3, 5$, and (b) $\mu B = 4$ and $\nu\tau_0 = 0.2$–2.0.

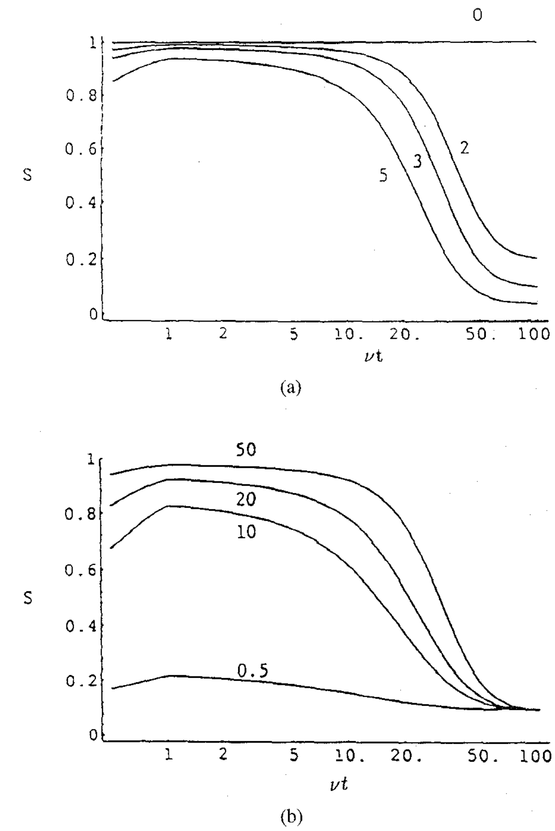


Fig. 2. Temporal evolution of the switch factor $S(t)$ for (a) $Q = 50$ and $\mu_0 B = 0$–5, and (b) $\mu_0 B = 3$, $Q = 0.5, 10, 20, 50$.

where n_0 is the final density or the asymptotic limit of $n(t)$ at late times. The form of (5) reveals that in addition to the conventional conduction current, represented by the first term in (5), there is a transient current flow, represented by the second term in (5). This is different from zero only when $Q(t) = 0$. If the pulse length satisfies $\tau_0 \gg 1/\nu$ this current becomes negligible and the conduction current dominates the response. In this case the presence of the magnetic field reduces the current amplitude by a factor $1 + \mu^2 B^2$. This is the well-known effect of magnetic insulation. For $\tau_0 \ll 1/\nu$ the transient current plays an important role. Its importance becomes apparent if we consider $Q(t) = n_0\delta(t)$. From (5) we find

$$\mathbf{J}(t) = \frac{\mathbf{J}_o}{1+\mu^2 B^2}[1 + e^{-\nu t}(\mu B \sin \Omega t - \cos \Omega t)] \quad (6a)$$

$$\mathbf{J}_0 \equiv n_0 e \mu \mathbf{E}. \quad (6b)$$

For $B = 0$ (6) becomes

$$J(t) = J_0(1 - \epsilon^{-\nu t}) \quad (7)$$

indicating a monotonic current response to the laser pulse, with time scale $1/\nu$. This is, however, not the case for $\mu B \gg 1$. It is clear that over time $1/\Omega$ the current overshoots its asymptotic value by a factor μB. This is equivalent to transient breakdown of the current with pulsewidth $t \lesssim 1/\nu$ driven in response to the short photonic pulse.

To summarize, this brief analysis of (6) revealed that the presence of an ambient magnetic field in a photoconducting switch results in a smaller value of the asymptotic current (magnetic insulation). Furthermore, for short pulses, the magnetic insulation is broken resulting in nonmonotonic current behavior over collisionless timescales. In the remainder of this paper, we examine numerically the behavior of (1)–(3) separately in the collisionless and collisional $(t \gg 1/\nu)$ regime. It is shown that in addition to the collisionless breakdown of magnetic insulation, similar behavior is exhibited in the collisional regime if the collision frequency ν is a strong function of the carrier density. A simple experiment confirming this behavior is also presented.

A. Collisionless Regime

The extent and limitations of the collisionless insulation breakdown were studied by numerically solving the set of (1)–(3) for an exponentially decaying source

$$Q(t) = (n_0/\tau_0)\exp(-t/\tau_0). \quad (8)$$

Fig. 1 shows the temporal response of the normalized current computed from (1)–(3) for different values of $\nu\tau_0$ and μB and for $\tau = 40\tau_0$, assuming zero initial carrier density. The current density $J(t)$ shown is normalized to the conduction current $J_c = J_0/(1+\mu^2 B^2)$. The results shown in Fig. 1(a) for $\nu\tau_0 = 0.2$ are identical to the response due to a delta-function

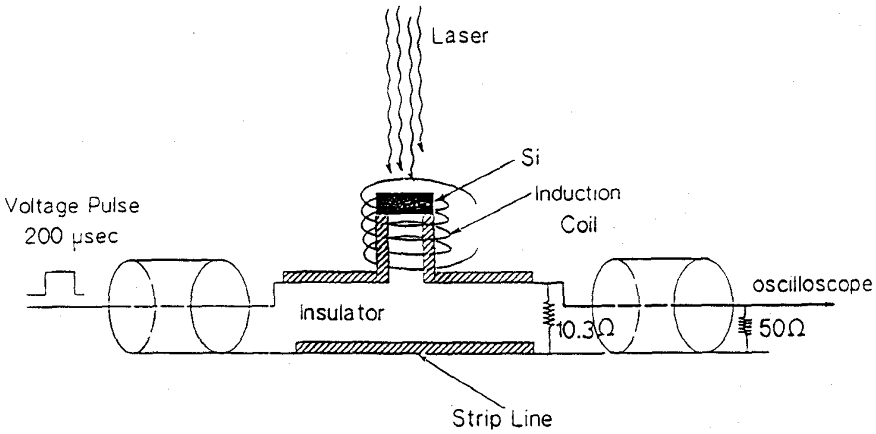


Fig. 3. Experimental setup.

pulse. The following points are apparent from Fig. 1(a) and (b):

1) For $\mu B < 1$ the current response to the laser pulse is monotonic independently of the value of τ_0. The response is similar to the unmagnetized case, and switching action cannot be achieved on a time scale shorter than the carrier loss time.
2) For $\mu B > 1$ and $\nu\tau_0 < 1$ the laser pulse induces a current pulse with time scale on the order of $1/\nu$. The amplitude of the pulse overshoot scales as μB, while the baseline conduction current scales as $1/(\mu B)^2$.
3) For long pulses $\nu\tau_0 \gg 1$, the velocity reduces to the constant term in (4). The current response again becomes monotonic, since in this case the maximum value of the transient current cannot exceed the value of the conduction current.

B. Collisional Regime

As noted previously, for pulses with $\nu\tau_0 \gg 1$ the transient current becomes negligible and the current response is dominated by the conduction current at all times. In this case nonmonotonic current response can be achieved only on time scales related to the loss time. It was shown by Dorquel and Lecturcq [7] that for energetic laser pulses the mobility becomes a function of the carrier density; as a result, the density dependence of the mobility should be included in the analysis. If the time scales are ordered according to $1/\nu,\ 1/\Omega \ll \tau_0 < \tau$ (so that the carrier density changes slowly compared to the other scales), the response to the laser pulse is given by the asymptotic form of (6a)

$$J(t) = en(t)\mu(n)/[1 + (\mu(n)B)^2] \tag{9}$$

with $n(t)$ given by the solution of (3) (see the Appendix). The functional form of $\mu(n)$ depends on the particular semiconductor. To demonstrate the role of the magnetization we assume here that

$$\mu(n) = \mu_0/[1 + (n/N)^{3/4}] \tag{10}$$

where N is a scale factor and μ_0 is the lattice mobility. It was shown in [7] that (10) accurately reflects the mobility dependence on carrier density for silicon. Using (9) and (10) and defining the normalized density as $x(t) = n(t)/N$ we find

$$J(t) = J_1\{x(t)/[1 + x(t)^{3/4}]\}/ \\ \{1 + [\mu_0 B/(1 + x(t)^{3/4})]^2\} \tag{11a}$$

$$J_1 = Ne\mu_0 E. \tag{11b}$$

The magnetic field effect in this case enters through the switch factor

$$S(x(t),\, \mu_0 B) = 1/[1 + (\mu_0 B/(1 + x(t)^{3/4})^2)]. \tag{12}$$

For $x(t) < 1$ the switch factor S has its asymptotic value equal to $1/[1 + (\mu_0 B)^2]$, which for large values of $\mu_0 B$ represents the effect of magnetic insulation. For $\mu_0 B \gg 1$ the current is reduced by a factor $\sim(\mu_0 B)^2$. On the other hand, for $x(t) \gg 1$ the switch factor has a value $S = 1$ and the current response is not affected by the magnetic field. Thus the magnetic insulation is broken.

The parametric dependence of the switching action on the magnetic field and laser pulse energy can be explored by considering a square laser pulse of duration τ_0, whose total number of photoconductivity-inducing photons per unit volume is P. In this case $x(t)$ will be given by

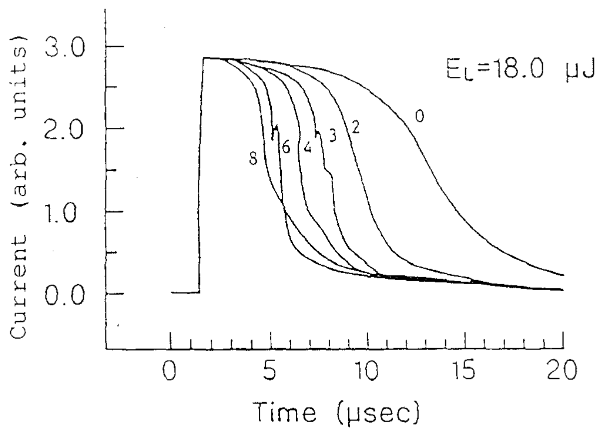
$$x(t) = Q(t), \quad t \le 1 \tag{13a}$$

$$x(t) = Q\exp[(1 - t)/\tau_1], \quad t > 1 \tag{13b}$$

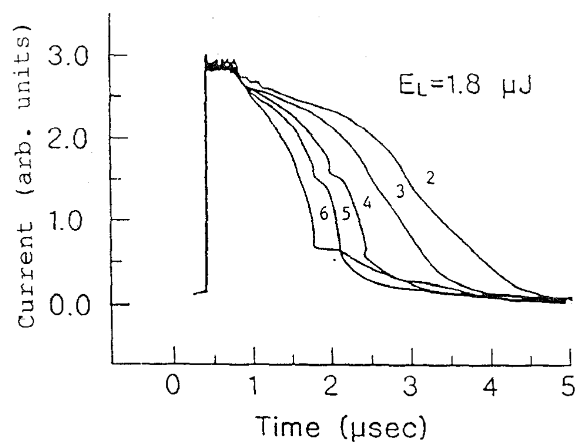
$$Q = P/N. \tag{13c}$$

In (13) the time is normalized to τ_0.

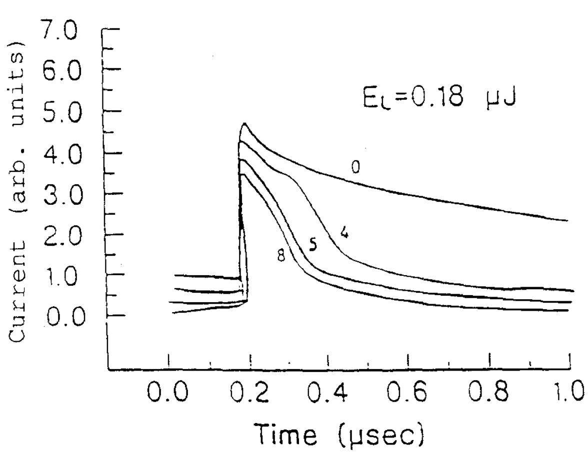
The system behavior is illustrated by solving (11) and (13) for different values of the parameters Q, $\mu_0 B$, and τ_1. Fig. 2(a) shows the temporal characteristics of $S(t)$ for $Q = 50$, $\tau_1/\tau_0 = 10$, $\nu\tau_0 = 200$, and $\mu_0 B = 0$–5. For $B = 0$ the value of the switch is $S = 1$ at all times. Notice that during the laser pulse, i.e., $t \le 1$, the insulation is essentially broken and the current is essentially equal to the unmagnetized current. However, for longer times, S decreases and reaches its asymptote of $1/(1 + \mu_0^2 B^2)$. As a result the current is significantly reduced during the off state. Fig. 2(b) indicates the role of the laser pulse energy represented by Q. It shows the results for $\mu_0 B = 3$ and $Q = 0.5$–50.



(a)



(b)



(c)

Fig. 4. Experimental results of the semiconductor current response to a laser pulse for (a) pulse energy 18 μJ, $B = 0, 2, 3, 4, 6, 8$ T, (b) pulse energy 1.8 μJ, $B = 2, 3, 4, 5, 6,$ T, and (c) pulse energy 0.18 μJ, $B = 0, 4, 5, 8$ T.

For $Q = 0.5$ there is no switching, but simply a uniform reduction of the unmagnetized current by a factor of 0.2 due to magnetic insulation. For $Q = 50$ there is complete breakdown of the magnetic insulation for $t \leq 1$, while 60–70% of the unmagnetized current is transmitted for $Q = 10$. On the other hand, the switching off is much faster for $Q = 10$ than $Q = 50$. These are indicative of the combination of factors that control the magnetized semiconductor response to the laser pulse in the long pulse limit.

Next we present a set of experimental results designed to test some of the theoretical concepts advanced above. Technical limitations allowed the study of phenomena induced by long pulses ($\tau_0 \ll 1/\nu$) only. A high-resistivity (20 kΩ/cm) silicon switch crystal with dimensions $3 \times 2 \times 0.25$ mm was placed between two identical strip transmission lines. The dimensions of each strip line were $15 \times 3 \times 0.2$ cm and its impedance was 1.6 Ω. The switch and part of the transmission line were mounted in the center of the induction coil as shown in Fig. 3. High magnetic fields up to 10 T were generated by discharging a 35-μF capacitor charged to a voltage of up to 6 kV into the induction coil. The magnetic field remained constant for times of about 50 μs, much longer than the time scales of interest.

The switch was illuminated with an 8-ns pulse from a Q-switched NdYAG laser (1.06 μm). The laser intensity incident on the switch varied from 3×10^2 W/cm^2 to 3×10^4 W/cm^2. The switch response time was measured for applied voltages between 10 and 500 V. A pulsed voltage source of 200 μs was used to prevent thermal runaway. The laser pulse was applied at the time of the maximum magnetic field value, causing photoconduction on the switch and turning on a current. The switch temporal response was measured for magnetic fields between 0 and 10 T.

A set of experiments performed with $B = 0$ for bias electric fields below 5 kV/cm indicated that the carrier loss time was inversely proportional to the bias voltage, as expected for sweep-out loss. The main results of the experimental investigations are summarized in Fig. 4. Fig. 4(a)–(c) shows the current response (expressed as the voltage drop across a load) for laser illumination of 18 μJ, 1.8 μJ, and 0.18 μJ, respectively, and for B = 0–8 T. In the first two cases the value of the maximum current and the characteristics of the current rise are independent of the magnetic field, as expected for large values of Q. On the other hand, the magnetic field exerts significant control over the current decay profile. The onset of the magnetic switch factor is clearly seen in Fig. 4. Shorter response times occur for higher values of the magnetic field, with properties consistent with the scaling of (11) and (13) and Fig. 2. The effect of the laser pulse energy can be seen by comparing Fig. 4(a), (b), and (c). Reduction of the laser energy results in faster decay times following the end of the laser pulse. However, the level breakdown of the magnetic insulation is also reduced, as seen in Fig. 4(c). Both trends are qualitatively consistent with the theoretical results shown in Fig. 2(b).

In this paper, we have presented a preliminary study of the effect of a strong magnetic field on the current response of a photoconductor driven by a laser pulse. The results are highlighted by:

- the effect of magnetic insulation, which significantly reduces the value of the dark and asymptotic currents;
- breakdown of the magnetic insulation due to polarization currents for $\nu\tau_0 < 1$, and by exploiting the dependence of the collision frequency on the carrier density for $\nu\tau_0 \gg 1$.

We have implicitly assumed that the two carrier species have comparable mobilities and carry comparable currents. Otherwise, a charge imbalance would develop, and an ambipolar electric field E_y would be present, whose magnitude is determined by the condition that the rates of charge loss of the two species are the same. This effect will be investigated elsewhere.

To our knowledge, these effects have not been discussed in the literature. It should be noted that the role of the magnetic field in photoconductors was previously studied experimentally by Moyer *et al.* [8] and theoretically by Parikh *et al.* [9]. In both cases, in contrast to the present case, only the regime $\mu B \ll 1$ was explored. Furthermore, the physical issues addressed concerned reduction of the carrier sweep-out time due to induced $E \times B$ carrier drift. In contrast, the mechanism described here is a bulk effect, which dominates in the limit of very strong magnetic field. In situations of practical interest both mechanisms may be important. Extension of our results to include multiple carriers is straightforward and will be published elsewhere. Before closing, we should emphasize the potential technological impact of the breakdown of the magnetic insulation due to polarization current. In this case the decay of the current occurs at times comparable to the cyclotron or the collision frequency *independently* of the carrier loss time and thereby without restrictions to the size of the switch. Notice that magnetic fields of 1 T give cyclotron frequencies of 1.3×10^{12} Hz for GaAs and 10^{13} Hz for InSb.

APPENDIX

Equation (1) and the solution (4) of the velocity equation presuppose that $\mathbf{E}$, $\mathbf{B}$, and ν are constant in time. If this is not true, the velocity of a carrier at time t depends on the time t' of its creation as well as its age, and the expression for the current becomes

$$\mathbf{J} = e\int_0^t dt'\dot{n}(t')\mathbf{v}(t,t'). \tag{A1}$$

The general solution of (2) for $\mathbf{v}(t',t') = 0$ is conveniently written in complex form as

$$w(t,t') = \frac{e}{m}\int_{t'}^{t} \mathcal{E}(t'') \times \exp\left[i\int_t^{t''} dt'''\{\nu(t''') + i\Omega(t'')\}\right] \tag{A2}$$

where $w = v_x + iv_y$ and $\mathcal{E} = E_x + iE_y$. Substituting (A2) in (A1), integrating by parts, and using $n(0) = 0 = w(t,t)$, we find

$$\mathcal{J}(t) = \frac{e^2}{m}\int_0^t dt' n(t')\mathcal{E}(t') \times \exp\left[i\int_t^{t'} dt''\{\nu(t'') + i\Omega(t'')\}\right] \tag{A3}$$

where $\mathcal{J} = J_x + iJ_y$.

In the present application $\mathcal{E}$ and Ω are constant, but ν depends on t through n by virtue of the assumed dependence (10). In the collisional limit, however, we can invoke the ordering specified prior to (9) in order to simplify this result. For this purpose we integrate by parts again, obtaining

$$\mathcal{J}(t) = \frac{e^2\mathcal{E}}{m}\left\{\left.\frac{n(t')\exp\left[i\int_t^{t'} dt''\{\nu(t') + i\Omega\}\right]}{\nu(t') + i\Omega}\right|_{t'=0}^{t'=t} - \int_0^t dt' \exp\left[i\int_t^{t'} dt''\{\nu(t'') + i\Omega\}\right] \times \frac{d}{dt'}\left[\frac{n(t')}{\nu(t') + i\Omega}\right]\right\}. \tag{A4}$$

The contribution from evaluating the integrated term at $t' = 0$ vanishes because $n(0) = 0$; but even if that were not the case, it would be exponentially small in consequence of the assumption $\nu t \gg 1$.

Integrating by parts repeatedly and discarding exponentially small terms at each stage, we find

$$\mathcal{J}(t) \approx \frac{e^2\mathcal{E}}{m}\left\{\frac{n(t)}{\nu(t) + i\Omega} - \frac{1}{\nu(t) + i\Omega}\frac{d}{dt}\frac{n(t)}{\nu(t) + i\Omega} + \frac{1}{\nu(t) + i\Omega}\frac{d}{dt}\left[\frac{1}{\nu(t) + i\Omega}\frac{d}{dt}\frac{n(t)}{\nu(t) + i\Omega}\right] \mp \cdots\right\}. \tag{A5}$$

It is clear that each successive term in the expansion is smaller than its predecessor by a quantity of order $\dot{\nu}/\nu^2 \sim 1/\nu t \ll 1$. The real part of the leading term is just (9).

ACKNOWLEDGMENT

Discussions with R. Shanny and P. Koert are gratefully acknowledged. The constructive criticism of an anonymous referee is greatly appreciated.

REFERENCES

[1] C. H. Lee, Ed., *Picosecond Optoelectronic Devices*. New York: Academic, 1984.
[2] C. H. Lee, "Optical control of semiconductor closing and opening switches," *IEEE Trans. Electron Devices*, vol. 37, pp. 2426–2437, 1990.
[3] W. C. Nunnally, "High-power microwave generation using optically activated semiconductor switches," *IEEE Trans. Electron Devices*, vol. 37, pp. 2439–2448, 1990.
[4] G. A. Mourou, C. V. Stancampiano, and D. Blumenthal, "Picosecond microwave pulse generation," *Appl. Phys. Lett.*, vol. 38, pp. 470–472, 1981.
[5] C. H. Lee, "Picosecond optics and microwave technology," *IEEE Trans. Microwave Theory Tech.*, vol. 38, pp. 596–607, 1990.
[6] D. H. Auston, "Subpicosecond electro-optic shock waves," *Appl. Phys. Lett.*, vol. 43, pp. 713–715, 1983.
[7] J. M. Dorkel and P. Leturcq, "Carrier mobilities in silicon semi-empirically related to temperature, doping, and injection level," *Solid-State Electron*, vol. 24, p. 821, 1981.
[8] R. H. Moyer, P. Agmon, T. L. Koch, and A. Yariv, "Ultrafast-magnetophotoconductivity of semi-insulating gallium arsenide," *Appl. Phys. Lett.*, vol. 39, pp. 266–268, 1981.
[9] C. D. Parikh and F. Lindholm, "Magnetic field effects in a photoconductive switch," *J. Appl. Phys.*, vol. 64, pp. 6546–6551, 1988.

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