Impact ionization fronts in semiconductors: superfast propagation due to “nonlocalized” preionization

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We discuss a new mode of ionization front passage in semiconductor structures. The front of avalanche ionization propagates into an intrinsic semiconductor with a constant electric field $E_m$ in presence of a small concentration of free nonequilibrium carriers - so called preionization. We show that if the profile of these initial carriers decays in the direction of the front propagation with a characteristic exponent $\lambda$, the front velocity is determined by $v_f \approx 2\beta_m/\lambda$, where $\beta_m \equiv \beta(E_m)$ is the corresponding ionization frequency. By a proper choice of the preionization profile one can achieve front velocities $v_f$ that exceed the saturated drift velocity $v_s$ by several orders of magnitude even in moderate electric fields. Our propagation mechanism differs from the one for well-known TRAPATT fronts. Finally, we discuss physical reasons for the appearance of preionization profiles with slow spatial decay.

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Propagation of impact ionization fronts in semiconductor structures represents a spectacular nonlinear effect which has important applications in pulse power electronics. In reverse-biased $p^+-n-n^+$ diode structures ionizing fronts propagate faster than the saturated drift velocity $v_s$. Such superfast propagation is possible due to the presence of small concentrations $n_0p_0$ of free electrons and holes in the depleted region. These free carriers which initiate an avalanche multiplication are often coined as “pre-ionization” of the medium. According to the conventional concept of ionization fronts in TRAPATT (TRAped Plasma Avalanche Triggered Transit) diodes the avalanche multiplication occurs within the ionization zone of length $\ell_f = \varepsilon_0(E_m - E_b)/qN_d$ where electric field exceeds the effective threshold of impact ionization $E_b$ (Fig. 1, curve 1). This length is finite due to the slope of the electric field in the $n$ base $dE/dx = qN_0/\varepsilon_0$ which depends on the doping level $N_d$ (note that $n_0p_0 \ll N_d$). The finiteness of the ionization zone $\ell_f$ prevents a uniform avalanche multiplication in the whole $n$ base and thus ensures the existence of the traveling front mode of avalanche breakdown. However, this concept is not applicable to $p-i-n$ structures with intrinsic ($N_d = 0$) base (Fig. 1, curve 2) as well as to short overvoltaged structures because in both cases $E > E_b$ in the whole $n$ base (Fig. 1, curve 3). On the basis of TRAPATT-like front concept one would expect that in these two cases pre-ionization of the high-field region triggers quasiform breakdown ruining the traveling front mode.

In this paper we argue that superfast impact ionization fronts are nevertheless possible in $p-i-n$ structures where $E > E_b$ everywhere in the high-field region providing the concentration profile of initial carriers $n_0(x), p_0(x)$ decays in the direction of front propagation. The propagation mechanism of such front is completely different from the conventional TRAPATT-like front. We find the front velocity analytically and show that it is controlled by the slope of pre-ionization profile, and that it can exceed $v_s$ by several orders of magnitude.

We consider a planar impact ionization front and describe it by the standard drift-diffusion model which consists of continuity equations for electron and holes concentrations $n$ and $p$ and the Poisson equation for the electric field $E$. For a self-similar front motion with constant velocity $v_f$ these equations can be simplified by introducing new variables $\sigma \equiv n + p$, $\rho \equiv p - n$. For intrinsic semiconductor ($N_d = 0$) the equations for $\sigma$, $\rho$ and $E$ in the co-moving frame $z = x + v_f t$ become

$$\frac{d}{dz} [v_f \sigma + v(E) \rho] - D \frac{d^2 \sigma}{dz^2} = 2 \alpha(E) v(E) \sigma, \quad (1)$$

$$v(E) \sigma + v_f \rho - \frac{D d\rho}{dz} = 0, \quad (2)$$

$$\frac{dE}{dz} = \frac{q}{\varepsilon_0} \rho, \quad (3)$$

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where \( v(E) \) is the drift velocity, \( D \) is the diffusion coefficient and \( \alpha(E) \) is the impact ionization coefficient. Here we neglect recombination and assume that electrons and holes are identical in a sense that \( \alpha(E) = \alpha_n(E) = \alpha_p(E) \) and \( v(E) = v_n(E) = v_p(E) \). For an infinite domain the boundary conditions are \( E \to E_m, \ \sigma, \ \rho \to 0 \) for \( z \to -\infty \) and \( E, \rho \to 0, \ \sigma \to \sigma_m \) for \( z \to +\infty \).

In the simplified case of \( D = 0 \) we use Eq. (2) to exclude \( \rho \) from Eqs. (1) and (3). This yields

\[
\frac{d}{dz} \left[ \frac{v_f^2 - v^2(E)}{v_f} \sigma \right] = 2 \alpha(E) v(E) \sigma, \tag{4}
\]

\[
\frac{dE}{dz} = -\frac{q}{\varepsilon \varepsilon_0} \frac{v(E)}{v_f} \sigma, \tag{5}
\]

\[
v(E) \sigma + v_f \rho = 0. \tag{6}
\]

Then by dividing equations (4) and (5) and integrating over \( E \) we immediately find the dependence \( \sigma(E) \) in the moving front:

\[
\sigma(E) = \frac{2 \varepsilon \varepsilon_0}{q} \frac{v_f^2 - v^2(E)}{v_f} \int_{E}^{E_m} \alpha(E) \ dE. \tag{7}
\]

The plasma concentration far behind the ionization zone

\[
\sigma_{pl} = \frac{2 \varepsilon \varepsilon_0}{q} \int_{0}^{E_m} \alpha(E) \ dE \tag{8}
\]

depends only on the electric field \( E_m \) (Fig. 2). The dependences \( p(E) \) and \( n(E) \) in the traveling front

\[
p(E), n(E) = \frac{\varepsilon \varepsilon_0}{q} \frac{v_f}{v_f \pm v(E)} \int_{E}^{E_m} \alpha(E) \ dE. \tag{9}
\]

follow directly from Eqs. (7) and are shown in Fig. 3. Remarkably, a traveling front solution exists for any \( v_f \geq v_s \). Within the approximation \( D = 0 \) the slowest solution corresponds to the shock front (discontinuous at \( E = E_m \)) that travels with a saturated drift velocity \( v_f = v_s \).

The above analysis does not allow to select a physically relevant solution and hence to find the actual front velocity \( v_f \). The selection problem remains in the case of \( D \neq 0 \). This is a general feature of fronts propagating into linearly unstable state [see Ref. 2 and references therein]. Ionizing fronts belong to this class since the state \( (E = E_m, \sigma = 0) \) is unstable: due to \( E_m > E_b \) any amount of free carriers leads to avalanche multiplication. It has also been suggested and confirmed by numerical simulations that in gases\(^{[10,11]}\) and semiconductors\(^{[12]}\) ionizing fronts are so called pulled fronts. For a pulled front, the dynamics in the part of the front where avalanche multiplication and screening are essentially nonlinear is subordinated to the linear dynamics of the front tip which fully determines the propagation velocity.\(^{[2]}\) The dynamics of the front tip is described by the linearized (near the state \( E = E_m, \sigma = 0 \)) version of equations (12) with constant coefficients \( v(E) = v_s \) and \( \alpha(E) = \alpha(E_m) = \alpha_m \)

\[
v_f \frac{d\sigma}{dz} + v_s \frac{d\rho}{dz} - D \frac{d^2 \sigma}{dz^2} = 2 v_s \alpha_m \sigma, \tag{10}
\]

\[
v_s \sigma + v_f \rho - D \frac{d\rho}{dz} = 0, \ \alpha_m \equiv \alpha(E_m). \tag{11}
\]
Here we take into account that in ionizing fields \( v(E) = v_s = \text{const} \). Solutions of these linear equations are exponential functions \( \sigma(z), \rho(z) \sim \exp(\lambda z) \), where the dispersion relation \( \nu_f(\lambda) \) remains to be found.

It is known from the theory of pulled fronts that the actual front velocity strongly depends on the type of initial conditions. All possible initial conditions split in two classes that lead to qualitatively different dynamics. **Localized** conditions correspond to initial profiles \( \sigma(x, t = 0) \) that are steeper than the profile \( \exp(\lambda^* x) \) with a certain characteristic exponent \( \lambda^* \): \( \sigma(x) < C \exp(\lambda^* x) \) for \( x \to -\infty \), where \( C \) is an arbitrary constant. In this case the front profile eventually becomes smoother and asymptotically reaches the profile \( \sigma(z) \sim \exp(\lambda^* z) \) at \( z \to -\infty \) that propagates with linear marginal stability velocity \( v^* = v_f(\lambda^*) \). Any initial profile that is strictly equal to zero for sufficiently small value of \( x \) also represents a localized initial condition. **Nonlocalized** initial conditions correspond to profiles \( \sigma(x, t = 0) \) with slow spatial decay that do not meet the above mentioned condition \( \sigma(x) < C \exp(\lambda^* x) \) for \( x \to -\infty \) and hence are smoother than \( \exp(\lambda^* x) \). In this case the front velocity is fully determined by \( \sigma(x, t = 0) \): for the initial profile \( \exp(\lambda_0 x) \) with \( \lambda_0 < \lambda^* \) the front velocity is given by the dispersion relation \( v_0 = v_f(\lambda_0) \).

The dispersion relation \( v_f(\lambda) \) follows from the characteristic equation of Eqs. \( (10) \) and \( (11) \)

\[
\ell^2 \lambda^3 - 2\ell \left( \frac{v_f}{v_s} \right) \lambda^2 + \left[ \left( \frac{v_f}{v_s} \right)^2 - 1 \right] \lambda - 2 \alpha_m \left( \frac{v_f}{v_s} \right) = 0, \ell \equiv \frac{D}{v_s} \tag{12}
\]

and is explicitly given by (see Fig. 4)

\[
\frac{v_f(\lambda)}{v_s} = \frac{\alpha_m}{\lambda} + \sqrt{1 + \left( \frac{\alpha_m}{\lambda} \right)^2} + \ell \lambda. \tag{13}
\]

The critical steepness \( \lambda^* \) and the velocity \( v^* \) correspond to the minimum point and are given by

\[
\frac{\lambda^*}{\alpha_m} = \sqrt{\frac{1}{\alpha_m \ell} \left( 1 - \frac{\alpha_m \ell}{2} + A \right)}, \tag{14}
\]

\[
\frac{v^*}{v_s} = \sqrt{1 + 5\alpha_m \ell - \left( \frac{\alpha_m \ell}{2} \right)^2 + (4 + \alpha_m \ell)A}, \tag{15}
\]

where \( A = \sqrt{(\alpha_m \ell)^2/4 + \alpha_m \ell}. \)

The right branch of the \( v_f(\lambda) (\lambda > \lambda^*) \) corresponds to fronts whose velocity increases with steepness \( \lambda \) due to diffusion. According to the concept of localized initial conditions these fronts are unstable: their steep profiles eventually relax to the profile with exponential tip \( \exp[\lambda^* z] \) that propagates with the velocity \( v^* \). The left branch of the \( v_f(\lambda) \) dependence (\( \lambda < \lambda^* \)) corresponds to stable fronts whose velocity decreases with \( \lambda \). These fronts correspond to nonlocalized initial conditions which are in the focus of our interest.

Characteristic values of the dimensionless parameter \( \alpha_m \ell \equiv \alpha_m D/v_s \) are 0.1 for Si and 1 for GaAs devices. As follows from Fig. 5, \( \lambda^* > \alpha_m \) in the relevant interval. Physically it means that the front propagating with linear marginal stability velocity \( v^* \) is so steep that the validity of the drift-diffusion approximation is questionable. This problem disappears for much smoother fronts that correspond to nonlocalized initial conditions (left branch in Fig. 4) if \( \lambda < \lambda^* \).

Eq. \( (13) \) leads to a simple relation \( v_f/v_s = 2\alpha_m/\lambda \) for the ionizing front velocity in case of preionization with decay exponent \( \lambda < \lambda^* \). For such fast fronts the effect of diffusion is negligible; in particular, Eq. \( (15) \) is fully applicable. Although \( v_f \) increase with \( \alpha_m \) and hence with electric field \( E_m \), it is the ratio \( \alpha_m/\lambda \) which actually counts. This ratio can be arbitrarily large resulting in front velocities that exceed the saturated drift velocity by many orders of magnitude. It means that a proper choice of slowly decaying preionization profile gives the possibility to achieve fast front propagation in even moderate (with respect to \( E_b \)) electric fields. However, the concentration of electron-hole plasma generated by the front passage increases with \( E_m \) (Fig. 2). Generally, the
electromagnetic limitation \( v_f < c \), where \( c \) is the velocity of light, may be important: for \( v_f \) comparable to \( c \) the full set of Maxwell equations shall substitute the Poisson equation in the model because the feedback from a nonstationary electromagnetic field created by the front passage on the front dynamics becomes essential.

Pre-ionization profiles with slow spatial decay can appear due to photoionization by photons from dense electron-hole plasma behind the front. In this case, \( \lambda^{-1} \) can be roughly identified as the light absorption length. This mechanism is efficient in direct-band materials and can be relevant for planar fronts in GaAs diode structures\(^4\) as well as finger-like streamers in direct-band bulk semiconductors.\(^5\) Another mechanism is related to field-enhanced ionization of deep centers in Si \( p^+-n^- \) structures used in pulse power applications.\(^6\) These high-voltage structures possess “hidden” deep levels – process-induced defects – with low recombination activity.\(^7\)

Preionization of the high-field space charge region can be due to field-enhanced ionization of these deep centers embedded in the \( n \) base.\(^8\) This ionization is more efficient near the \( p^+-n \) junction where the electric field is stronger. Hence the profile of initial carriers decreases along the \( n \) base. The characteristic decay length \( \lambda^{-1} \) is expected to be a fraction of the \( n \) base width \( W \sim 100 \mu m \). At the same time for low doped \( n \) base the electric field can be above the ionization threshold \( E_b \) everywhere. This may result in front velocities that exceed \( v_s \) by several orders of magnitude. Numerical simulations of such triggering process will be presented elsewhere.

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