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Plasma Instability and Terahertz Generation in HEMTs Due to Electron Transit-Time Effect

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SUMMARY We study the coupled spatio-temporal variations of the electron density and the electric field (electron plasma oscillations) in high-electron mobility transistors using the developed device model. The excitation of electron plasma oscillations in the terahertz range of frequencies might lead to the emission of terahertz radiation. In the framework of the model developed, we calculate the resonant plasma frequencies and find the conditions for the plasma oscillations self-excitation (plasma instability) We show that the transit-time effect in the high-electric field region near the drain edge of the channel of high-electron mobility transistors can cause the self-excitation of the plasma oscillations. It is shown that the self-excitation of plasma oscillations is possible when the ratio of the electron velocity in the high field region, u_d, and the gate length, L_g, i.e., the inverse transit time are sufficiently large in comparison with the electron collision frequency in the gated channel, ν. The transit-time mechanism of plasma instability under consideration can superimpose on the Dyakonov-Shur mechanism predicted previously strongly affecting the conditions of the instability and, hence, terahertz emission. The instability mechanism under consideration might shed light on the origin of terahertz emission from high electron mobility transistors observed in recent experiments.

key words: *heterostructure, high-electron mobility transistor, plasma oscillations, plasma instability, transit-time effect, terahertz radiation*

1. Introduction

Heterostructure electron channels with high electron mobility can serve as resonant cavities for terahertz (THz) plasma waves. These waves consist in the self-consistent spatio-temporal variations of the electron density in the channel and the electric field around it [1], [2]. Plasma effects in the gated electron channel of heterostructures like high-electron mobility transistors (HEMTs) have attracted much attention. This is mainly due to promising prospects to use these effects for detection, frequency multiplication, and generation of THz radiation. Such effects have been extensively studied both theoretically and experimentally. In particular, THz emission by plasma oscillations in InGaAs and GaN HEMTs was recently reported by Knap et al. [3] and Deng et al. [4]. The self-excitation of plasma oscillations in and THz emission from the HEMTs in question was linked by the authors to an instability of a steady-state electron flow in the transistor channel associated with the reflection of the plasma waves from the drain edge of the gated channel, i.e., with the so-called Dyakonov-Shur (DS) mechanism of plasma wave instability [5]. The pertinent instability condition is given by $u_g/L_g > \nu/2$, where u_g, L_g, and ν are the effective electron drift velocity, the length of the gate, and the frequency of electron collisions with impurities and phonons (inverse momentum relaxation time) in the HEMT gated channel, respectively. As shown in Ref. [4], the plasma wave instability in GaN HEMTs took place in fairly long channel devices, where this condition was not met. Hence, more complex mechanism of the plasma wave self-excitation can be responsible for the plasma instability in HEMTs. We showed previously [6]–[8] that transit time (TT) effects related to the electron current flowing to the gate contact might lead to the negative dynamic resistance and, therefore, greatly enhance the plasma wave growth increment in different heterostructure devices. Later on, we observed the peak of the device response in the regime when the gate leakage current was large [9], which agreed with the predictions of the theory [6], [7].

In this paper, we study the plasma oscillations in HEMTs and show that the TT effect related to the electron drift across the high field region that appears in the drain side of the channel at voltages exceeding the HEMT saturation voltage can greatly enhance the increment of the plasma wave instability. This mechanism might contribute to the plasma oscillations leading to THz emission reported recently [3], [4], since this emission occurs at voltages close to or above the drain saturation voltage. The present paper presents a substantial extension and generalization of our recent letter [10].

2. Equations of the Model

The mechanism under consideration is associated with the following two features of HEMTs: (1) the electron system in the HEMT channel under the gate can serve as a resonant cavity (due to existence of plasma resonances) and (2) the electron transit delay in the channel portion between the gate and the drain can result in negative dynamic conductivity at the frequencies determined by the inverse time of the electron transit.

A schematic view of the n-type HEMT structure under consideration is shown in Fig. 1. We describe the dynamics of electrons using the hydrodynamic electron transport model which comprises the Euler equation and the continuity equation for the electron system and the Poisson equation for the self-consistent electric potential [1], [2], [5]. The lin-

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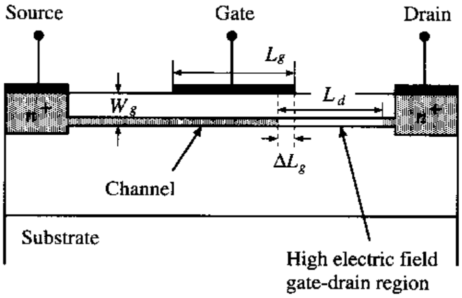

Fig. 1 Schematic view of the HEMT structure.

earized versions of equations of the model for the ac components (with the signal frequency ω) of the electron hydrodynamic velocity $u_\omega(x)\exp(-i\omega t)$, the electron sheet concentration $\Sigma_\omega(x)\exp(-i\omega t)$ and the self-consistent potential $\varphi_\omega(x)\exp(-i\omega t)$ in gated portion of the channel can be presented as

$$(\nu - i\omega)u_\omega + u_\omega\frac{du_g}{dx} + u_g\frac{du_\omega}{dx} = \frac{e}{m}\frac{d\varphi_\omega}{dx}, \quad (1)$$

$$-i\omega\Sigma_\omega + \Sigma_0\frac{du_\omega}{dx} + u_\omega\frac{d\Sigma_g}{dx} + u_g\frac{d\Sigma_\omega}{dx} + \Sigma_\omega\frac{du_g}{dx} = 0, \quad (2)$$

$$\frac{\varphi_\omega}{W_g} = -\frac{4\pi e}{\text{æ}}\Sigma_\omega. \quad (3)$$

Here $\Sigma_g(x)$ and $u_g(x)$ are the dc electron concentration and drift velocity in the channel W_g is the thickness of the gate layer separating the channel from the gate, $e = |e|$ and m are the electron charge and effective mass, and æ is the dielectric constant. The quantities $\Sigma_g(x)$ and $u_g(x)$ are determined by the donor concentration in the gate layer and the channel, Σ_d, by the conduction band discontinuity at the heterointerface, by the gate and the source-drain voltages V_g and V_{sd}, and, in the case of GaN-based HEMTs, by the polarization field. The axis x is directed along the HEMT channel. Equation (3) corresponding to the gradual channel approximation [11] can be used when $W_g \ll L_g$. The dc current along the channel is equal to $J_0 = e\Sigma_g(x)u_g(x) = const$. To treat the problem analytically, we assume that the electron concentration in the gated portion of the channel $\Sigma_g(x)$ to be independent of the coordinate. Hence $u_g(x)$ is also close to a constant. As a result, one can neglect the derivatives $d\Sigma_g(x)/dx$ and $du_g(x)/dx$ in Eqs. (1) and (2), so that Eqs. (1)–(3) are reduced to the following equation for the ac potential in the gated portion of the channel:

$$\frac{d^2\varphi_\omega}{dx^2} + \frac{iu_g(2\omega + i\nu)}{(s^2 - u_g^2)}\frac{d\varphi_\omega}{dx} + \frac{\omega(\omega + i\nu)}{(s^2 - u_g^2)}\varphi_\omega = 0, \quad (4)$$

where $s = \sqrt{4\pi e^2\Sigma_g W/\text{æ}m}$ is the characteristic plasma wave velocity [5]–[8]. Since Σ_g depends on the gate voltage V_g, the plasma wave velocity and, consequently, the plasma frequencies can be tuned by this voltage.

When the length of the ungated section of the channel, L_s, between the source and gate is smaller than L_g, this section simply contributes to the device series resistance weakly affecting the plasma oscillations in the HEMT structure [12]. Neglecting the potential drop across the ungated section of the channel between the source and gate, one can set

$$\varphi_\omega|_{x=0} = 0. \quad (5)$$

The ac current at the drain edge of the channel, i.e., at $x = L_g$, is equal to $J_\omega|_{x=L_g} = e[u_g\Sigma_\omega + \Sigma_g u_\omega]|_{x=L_g}$. In the saturation region, the distribution of the electron velocity in the HEMT channel could be approximated by dividing the channel into two regions [10]: the gated section of the channel ($x \leq L_g^{(eff)}$, Region I), where the electron velocity, u_g, is low and the electron concentration is high) and the high-field gate-drain "velocity saturation region" ($L_g^{(eff)} < x < L_g^{(eff)} + L_d$, Region II), where the electron velocity, u_d, as well as the longitudinal electric field are high, and the electron concentration is low. Here $L_g^{(eff)} \leq L_g$ ($L_g^{(eff)} = L_g - \Delta L_g$, see Fig. 1). The feasibility of such a partition is confirmed by the Monte Carlo calculations [13], [14] showing, in particular, that in short channel devices, u_d can significantly exceed the electron saturation velocity.

Considering this two-region approximation neglecting for simplicity the distinction between $L_g^{(eff)}$ and L_g (see, however, a comment in Sec. 5), the electron current injected into the high-electric field gate-drain region (at $x = L_g$) can be presented as $J_\omega^{inj}|_{x=L_g+0} = eu_d\Sigma_\omega|_{x=L_g} = -(\text{æ}u_d/4\pi W_g)\varphi_\omega|_{x=L_g}$. The electrons propagating in the high-field gate-drain region induce the ac current in the gated channel and the drain contact. Assuming that electrons in propagate with constant velocity, u_d, one can find that the electron ac concentration as a function of the coordinate varies as $\Sigma_\omega(x) = \Sigma_\omega|_{x=L_g}\exp[i\omega(x - L_g)/u_d]$. As a result, following the Shockley-Ramo theorem [15], [16], the ac current induced in the gated channel can be presented as

$$J_\omega = \frac{eu_d\Sigma_\omega|_{x=L_g}}{L_d}\int_{L_g}^{L_g+L_d} dx g(x - L_g)\exp\left[i\frac{\omega(x - L_g)}{u_d}\right]$$
$$= -\frac{\text{æ}u_d\varphi_\omega|_{x=L_g}}{4\pi eW_gL_d}\int_0^{L_d} dx g(x)\exp\left(i\frac{\omega x}{u_d}\right). \quad (6)$$

Here $g(x)$ is the form-factor depending on the geometry of the conducting regions (quasi-neutral sections of the channel, gate, and drain contact). Its particular shape is determined by the dependence of the charge induced in highly conducting regions by the moving electron on the position of the latter. In the case when the induced ac current is associated primarily with the variation of the charges in the quasi-neutral sections of the channel (i.e., in lateral conducting regions, so that the gate-drain region can be considered as some kind of a slot diode), one obtains [17], [18] $g(x) = 2/\pi\sqrt{1 - (2x - L_d)^2/L_d^2} = g_l(x)$. If the quasi-neutral section of the channel immediately adjacent to a bulk drain contact is sufficiently short, the propagating electrons induce the ac current directly in this contact. In this case, $g(x) = 2/\pi\sqrt{1 - (x - L_d)^2/L_d^2} = g_b(x)$. Since the geometry of the conducting regions in HEMTs under applied voltage is fairly complex, we shall use the above form-factors bearing in mind that they represents limiting cases. Equalizing the electron ac current in the gated channel near its drain edge and the ac current induced in this channel by the electrons propagating between the gate and the drain, we arrive at the following condition:

$$\left.\frac{d\varphi_\omega}{dx}\right|_{x=L_g} = i\,F(\omega)\,\varphi_\omega|_{x=L_g}. \tag{7}$$

Here

$$F(\omega) = \frac{[-u_g(2\omega + i\nu) + u_d(\omega + i\nu)f(\omega)]}{(s^2 - u_g^2)} \tag{8}$$

with

$$f(\omega) = \frac{2\,r}{\pi L_d}\int_0^{L_d} dx \frac{\exp(i\omega x/u_d)}{\sqrt{1-(2x-L_d)^2/L_d^2}} = r\,J_0(\omega\tau_d/2)[\cos(\omega\tau_d/2) + i\sin(\omega\tau_d/2)], \tag{9}$$

and

$$f(\omega) = \frac{2\,r}{\pi L_d}\int_0^{L_d} dx \frac{\exp(i\omega x/u_d)}{\sqrt{1-(x-L_d)^2/L_d^2}} = r\,[J_0(\omega\tau_d) + iH_0(\omega\tau_d)] \tag{10}$$

for the abovementioned form-factors $g_l(x)$ and $g_b(x)$, respectively. Here, $J_0(x)$ and $H_0(x)$ are the Bessel and Struve functions [19], and $\tau_d = L_d/u_d$ is the time of electron transit between the edge of the gate and the drain. If the electron drift velocity in the gate-drain region markedly varies with the coordinate due to the velocity overshoot effect, τ_d is determined by some average drift velocity. The dependences of $\mathrm{Re}f(\omega)$ given by Eqs. (9) and (10) are qualitatively very similar. This, as will be seen in the following, means that the specific shape of the conducting regions is not an essential factor affecting the frequency dependence of the induced ac current and, therefore, the plasma oscillations and the conditions of plasma instability. However, to take into account a "parasitic" effect associated with some leakage of the ac current say, to the gate, we have introduced in Eqs. (9) and (10) a phenomenological factor r ($r \leq 1$). If the ac current induced by the electrons propagating in the high-field gate-drain region is neglected ($f(\omega) \to 0$), Eqs. (7)–(10) coincide with the pertinent equations obtained for the case of DS instability.

3. Dispersion Equation

Solving Eq. (4) with boundary conditions (5) and assuming for simplicity $\nu < |\omega|$ (and omitting, therefore, the terms of the order of $(\nu/\omega)^2$),we obtain

$$\varphi_\omega \propto \exp\left[-i\left(\frac{\omega + i\nu/2}{s^2 - u_g^2}\right)u_g\,x\right]\left\{\exp\left[i\left(\frac{\omega + i\nu/2}{s^2 - u_g^2}\right)s\,x\right] - \exp\left[-i\left(\frac{\omega + i\nu/2}{s^2 - u_g^2}\right)sx\right]\right\}. \tag{11}$$

After that, applying boundary condition (7) and taking into account that for real HEMTs, the values $(\nu/2\Omega)$, (u_g/s), and $(u_d/s)\cdot f(\omega)$ are actually small parameters ($|f(\omega)| = |J_0(\omega\tau_d/2)| < 1$), we arrive at the following equation which in fact is the dispersion equation for the plasma oscillations:

$$\cot\left(\frac{\pi}{2}\frac{\omega}{\Omega}\right) \simeq i\left[\frac{\left(\frac{\pi}{4}\frac{\nu}{\Omega}\right)}{\sin^2\left(\frac{\pi}{2}\frac{\omega}{\Omega}\right)} - \frac{u_g}{s} + \frac{u_d}{s}f(\omega)\right]. \tag{12}$$

Here $\Omega = \pi(s^2 - u_g^2)/2sL_g$ is the fundamental plasma frequency. Dispersion equation (12) differs in fact from that obtained previously [5] (which describes the DS instability) by the last term in the right-hand side associated with the electron TT effect. Considering the right-hand side of Eq. (12) as a perturbation and using Eq. (9), one can obtain the following formulas for the real and imaginary parts of the oscillation frequency:

$$\mathrm{Re}\,(\omega) = \omega_n \simeq (2n-1)\Omega, \tag{13}$$

$$\mathrm{Im}\,(\omega) \simeq -\frac{\nu}{2} + \gamma_{DS} + \gamma_{TT}. \tag{14}$$

Here,

$$\gamma_{DS} \simeq \frac{u_g}{L_g}, \tag{15}$$

$$\gamma_{TT} \simeq -\frac{u_d}{L_g}r\cos(\omega_n\tau_d/2)\,J_0(\omega_n\tau_d/2), \tag{16}$$

and $n = 1, 2, 3, \ldots$ is the plasma mode index. The quantity $\mathrm{Im}\,(\omega)$ is the increment (growth rate) of the amplitude of the plasma oscillation with the frequency $\mathrm{Re}\,(\omega)$. The case $\gamma_{DS} > 0$ corresponds to the contribution of the DS mechanism to the increment (growth rate) of the plasma oscillations, whereas γ_{TT} (which can be both positive and negative depending on $\omega_n\tau_d$) corresponds to the contribution of the electron TT mechanism. Calculating γ_{TT} considering Eq. (10) instead of Eq. (9), one can obtain

$$\gamma_{TT} \simeq -\frac{u_d}{L_g}r\,J_0(\omega_n\tau_d). \tag{17}$$

In GaAs-based HEMTs with sufficiently short gates, the fundamental plasma frequency Ω falls into the THz range of frequencies. Indeed, assuming the gate layer thickness W = 20 nm and and the dc electron concentration in the channel Σ_g = (3–8) $\times$ 10^{11} cm^{-2}, we find $s \simeq (5.0\text{–}8.0)\times 10^7$ c/ms. Setting $s = (5.0\text{–}8.0)\times 10^7$ c/ms and L_g = 100–250 nm, at $u_g \ll s$ one can obtain $\Omega/2\pi \simeq$ 0.5–2.0 THz.

4. Instability Conditions

The plasma instability occurs when the increment (plasma oscillations growth rate) becomes positive:

$$\mathrm{Im}\,(\omega) > 0. \tag{18}$$

Considering this and using Eqs. (14)–(16), one can arrive at the following criterion for the self-excitation of the first plasma mode ($n = 1$):

$$\frac{u_g}{L_g} - \frac{u_d}{L_g} r \cos\left(\frac{\pi}{4}\frac{s}{u_d}\frac{L_d}{L_g}\right) J_0\left(\frac{\pi}{4}\frac{s}{u_d}\frac{L_d}{L_g}\right) > \frac{\nu}{2}. \tag{19}$$

Since usually $s \gg u_d, u_g$, the transit angle Θ is not small. At $\Theta = (\pi/2)(s/u_d)(L_d/L_g) \gg 1$, inequality (19) reduces to

$$\frac{u_g}{L_g} > \frac{u_d}{L_g}\frac{2r}{\pi}\sqrt{\frac{u_d}{s}\frac{L_d}{L_g}} + \frac{\nu}{2} \simeq \frac{\nu}{2}. \tag{20}$$

In this case, the electron TT effect is insignificant and, hence the instability condition coincides with that related to the DS mechanism. This situation might correspond to a very large source-drain voltage (large L_d) and to a very large value of characteristic plasma velocity, s (large gate voltage swing).

At the transit angles corresponding to the commensurability of the plasma resonant frequency and the TT frequency, the second term in the left-hand side of Eq. (19) reaches maxima at which this term is positive. The situations when the fundamental plasma frequency ($n = 1$) is equal to the lowest two TT frequencies correspond to $\Theta = \Theta_1 \simeq 3.94 \simeq 1.25\pi$ and $\Theta = \Theta_2 \simeq 10.22 \simeq 3.25\pi$, so that $\cos\Theta_1 J_0(\Theta_1) \simeq -0.094$ and $\cos(\Theta_1/2) J_0(\Theta_1/2) \simeq -0.055$. Taking this into account and assuming $r = 1$, criterion (18) yields the following necessary condition for the instability at the first and second TT resonances:

$$\frac{u_g + 0.094u_d}{L_g} > \frac{\nu}{2}, \qquad \frac{u_g + 0.055u_d}{L_g} > \frac{\nu}{2}. \tag{21}$$

In this case, both mechanisms promote the plasma instability. If the ratio u_d/u_g is sufficiently large, the electron TT effect can be the main reason for the plasma instability.

If the transit-time increment is given by Eq. (17), one obtains the following instability criterion replacing inequality (19):

$$\frac{u_g}{L_g} - \frac{u_d}{L_g} r J_0\left(\frac{\pi}{2}\frac{s}{u_d}\frac{L_d}{L_g}\right) > \frac{\nu}{2}. \tag{22}$$

When $\Theta = \Theta_1 = 3.88 \simeq 1.235\pi$, $\Theta = \Theta_2 = 10.17 \simeq 3.24\pi$, and $r = 1$, inequality (22) yields

$$\frac{u_g + 0.40u_d}{L_g} > \frac{\nu}{2}, \qquad \frac{u_g + 0.25u_d}{L_g} > \frac{\nu}{2}. \tag{23}$$

Comparing conditions (21) and (23), one can see that the second one (corresponding to the situation when the bulk drain contact provides a good collection of the induced ac current) is more liberal. Assuming that $\nu = 5 \times 10^{11}$ s^{-1} (for materials like GaAs this electron collision frequency corresponds to the electron mobility in the HEMT channel about 150,000 cm^2/V s) $u_g = 1 \times 10^7$ c/ms, and $u_d = (2\text{–}4) \times 10^7$ c/ms, the plasma instability associated with both DS and TT mechanisms can, as follows from inequality (24), occur in HEMTs with $L_g <$ 320–1040 nm. As shown by recent Monte Carlo calculations [14], the electron drift velocity in the high-field gate-drain region in GaInAs HEMTs can

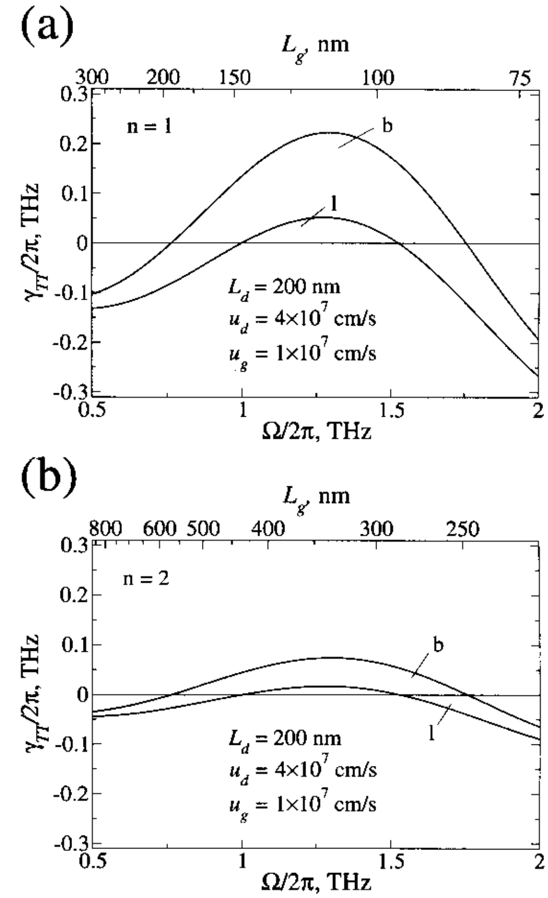

Fig. 2 Increment associated with the TT mechanism, γ_{TT}, versus fundamental plasma frequency Ω and gate length L_g in HEMTs with (l) lateral and (b) bulk configurations of the drain and characteristic plasma wave velocity $s = 5.8 \times 10^7$ c/ms ($\Omega \propto s/L_g$): (a) fundamental plasma mode ($n = 1$) and (b) second plasma mode ($n = 2$).

be as high as $u_d = 8 \times 10^7$ c/ms. At $\nu = 5 \times 10^{11}$ s^{-1}, $u_g = 1 \times 10^7$ c/ms, and $u_d = 8 \times 10^7$ c/ms, one can find that the first condition (24) satisfies if $L_g < 1680$ nm. Thus, the plasma instability associated with the TT effect can, in principle, occur in HEMTs with fairly long gates. However, in HEMTs with long gates, the combined resonance can be achieved only for higher plasma modes ($n > 1$) because the fundamental plasma frequency markedly decreases with increasing gate length ($\Omega \propto L_g^{-1}$).

Figure 2 demonstrates a quantitative difference in the dependence of the TT increments, γ_{TT}, in HEMTs with different form-factors, $g_l(x)$ and $g_b(x)$, respectively, calculated as a function of the fundamental plasma frequency Ω and gate length L_g (related to each other as ($\Omega \propto s/L_g$, where s is assumed to be fixed). using Eqs. (16) and (17). It is assumed that $s = 5.8 \times 10^7$ c/ms and $u_d = 4 \times 10^7$ c/ms. As seen from Fig. 2, despite qualitatively similar dependence on Ω and L_g, the TT mechanism in HEMTs with bulk drain contact is more pronounced. Comparison of Figs. 2(a) and 2(b) shows that the second plasma mode ($n = 2$) with $\text{Re}(\omega) = \omega_2 \simeq 3\Omega$) can excite in HEMTs with longer gate length, although its increment is smaller that that of the fun-

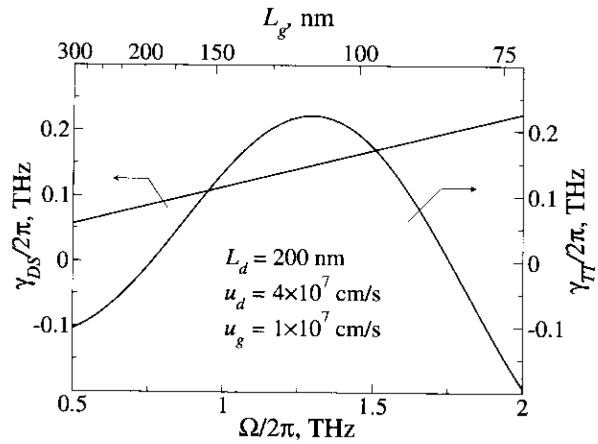

Fig. 3 Contributions to the plasma instability increment of the DS and TT mechanisms, ($\gamma_{DS}/2\pi$ and $\gamma_{TT}/2\pi$, versus fundamental plasma frequency Ω and gate length L_g ($s = 5.8 \times 10^7$ c/ms).

damental plasma mode ($n = 1$ and $\mathrm{Re}(\omega) = \omega_1 \simeq \Omega$).

Figure 3 shows the contributions of the pertinent mechanisms to the net increment of the plasma instability, $\mathrm{Im}\,(\omega)$, as a function of the fundamental plasma frequency Ω and gate length L_g (related to each other as ($\Omega \propto s/L_g$) calculated using Eqs. (15) and (17) for HEMTs with $\nu = 5 \times 10^{11}\ \mathrm{s}^{-1}$, $s = 5.8 \times 10^7$ c/ms and $L_d = 200$ nm at $u_g = 1 \times 10^7$ c/ms and $u_d = 4 \times 10^7$ c/ms. One can see from Fig. 3 that at relatively small values of Ω, the TT effect provides a negative contribution to the increment and can overcome the DS mechanism. For the sake of definiteness, in Fig. 3 as well as in Figs. 4 and 5, we show the results of calculations related to the fundamental plasma mode ($n = 1$).

Figure 4 shows the plasma wave instability increment $\mathrm{Im}\,(\omega)/2\pi$ as a function of the fundamental plasma frequency Ω and gate length L_g at fixed characteristic plasma wave velocity s calculated using Eqs. (14), (15), and (17) for different lengths of the high-field gate-drain region and different electron velocities in this region. At modest ratios u_d/u_g as in Fig. 4(a) where $u_d/u_g = 2$, the increment is positive in a wide ranges of Ω and L_g that corresponds to the plasma instability. In this case, the contribution of the TT mechanism reveals in an oscillatory dependence of $\mathrm{Im}\,\omega$ versus Ω and L_g. However at larger ratios u_d/u_g, the dependence of the increment on Ω and L_g, as shown in Figs. 4(b) and 4(c), is determined primarily by the TT effect.

Since the instability associated with the TT effect can occur when in some ranges of the transit angle $\Theta \propto L_d/L_g$, to achieve the plasma oscillations self-excitation, both L_d and L_g should be chosen properly. Figure 5 shows the instability region (where the increment is positive, $\mathrm{Im}\,\omega > 0$) on the L_d–L_g-plane.

5. Discussion

The suppression of the DS mechanism of instability at small values of $\Theta \propto (s/u_d)(L_d/L_g) \propto \tau_d$ can be interpreted as follows. The DS mechanism is associated with the reflection

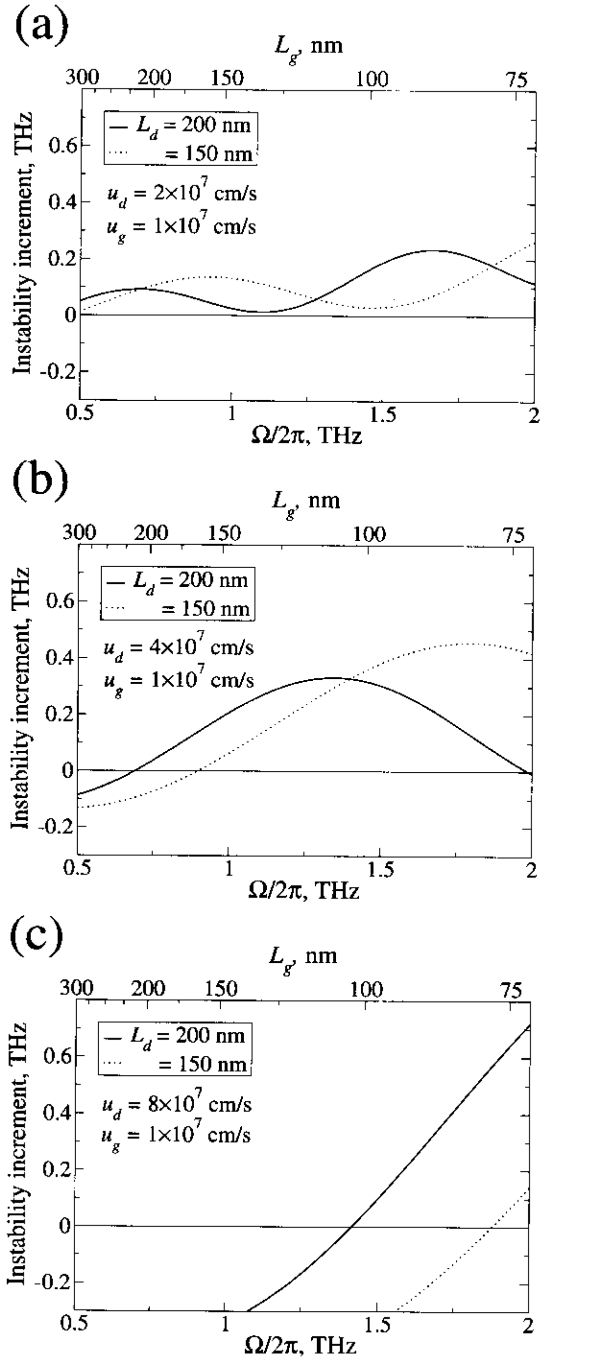

Fig. 4 Plasma instability increment $\mathrm{Im}\,(\omega)/2\pi$ versus fundamental plasma frequency Ω and gate length L_g for different lengths of high-field gate-drain region L_d ($L_d = 200$ nm and $L_d = 150$ nm) and different electron velocities u_d in this region: (a) $u_d = 2 \times 10^7$ c/ms, (b) $u_d = 4 \times 10^7$ c/ms, and (c) $u_d = 8 \times 10^7$ c/ms with $s = 5.8 \times 10^7$ c/ms.

of the plasma waves from the drain edge of the gated channel. Such a reflection is effective if the conductivity of the high-field gate-drain region at the frequencies under consideration is sufficiently small compared with the conductivity of the gated section of the channel. This case corresponds to Eq. (7) with $f(\omega) = 0$, so that $F(\omega) \propto -u_g\omega/s^2 < 0$. However, at large source-drain voltages, V_{sd}, the low-frequency

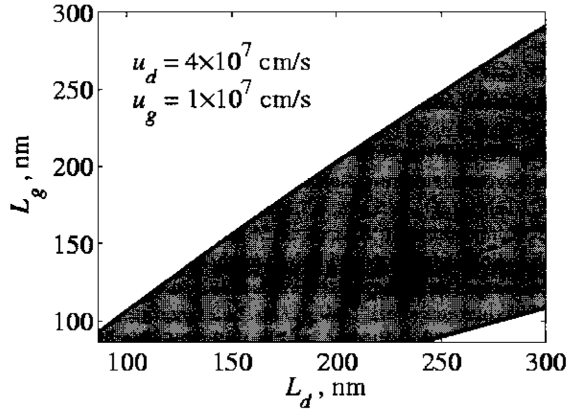

Fig. 5 Plasma instability region (shaded) corresponding to the positive increment, $\mathrm{Im}\,\omega > 0$, on the L_d-L_g-plane for $s = 5.8 \times 10^7$ c/ms. Dashed line corresponds to maxima values of increment $\mathrm{Im}\,\omega$.

conductivity of this region might be not that small due to relatively high values of u_d if L_d is moderate. This corresponds to short τ_d and, consequently, small Θ, at which $f(\omega) \simeq f(0) = r$. In such a case, $F(\omega) > 0$ (if, certainly, $u_g \ll u_d$ and $r \sim 1$). This means that the plasma waves are effectively "absorbed" in the high-field gate-drain region and, hence, their reflection is insufficient to provide necessary positive feedback for the occurrence of the plasma instability. At a large Θ, the conductivity of the high-field gate-drain region is rather small ($|f(\omega)| \ll 1$). As a result, $F(\omega) < 0$ that corresponds to an effective reflection of the plasma waves from the drain edge of the gated channel (positive feedback) which leads to the DS instability if the collision damping of the plasma waves is overcome (see condition (20)). The most interesting situation arises when Θ corresponds to a negative dynamic conductivity of the high-field gate-drain region in some range of frequencies associated with the electron TT effect. This occurs at the frequencies at which $f(\omega) < 0$. In this situation, the negative conductivity provides the positive feedback resulting in the instability. Indeed, if $f(\omega) < 0$, $F(\omega) \propto -[u_g - u_d f(\omega)]\omega/s^2 = -[u_g + u_d|f(\omega)|]\omega/s^2 < 0$. In this situation, inequalities (21) and (23) are satisfied.

The transit angle depends actually on the gate and drain voltages. This due to the dependence of the plasma wave velocity s on the gate voltage V_g and the dependence of the electron transit time τ_d on the source-drain voltage V_{sd}. It is natural to assume that τ_d and, hence, Θ decrease with increasing drain voltage, so that Θ can varies from rather large values at moderate V_{sd} to to those corresponding to the electron transit resonances. This might explain the occurrence of the plasma instability and THz emission when V_{sd} exceeds some threshold value (that was observed experimentally) when the contribution of the TT mechanism becomes pronounced. The disappearance of the instability and THz emission at fairly high drain voltages can be attributed, in such a case, to a decrease of the transit angle to the values at which the TT mechanism stabilizes the electron system. This is in qualitative agreement with the experimental data [3].

To account for the difference between $L_g^{(eff)}$ and L_g, i.e., the gate length modulation effect at drain voltages, V_{sd}, exceeding the source-drain saturation voltage, $V_{sd}^{(sat)}$, the value of L_g in the above equations has to be replaced by $L_g^{(eff)} = L_g - \Delta L_g$. Here ΔL_g can be crudely estimated as $(V_{sd} - V_{sd}^{(sat)})/E_m$, where E_m is the characteristic gate length modulation field. An increase in the drain bias can result in both an increase in the length of the gate-drain high-electric-field region L_d and pertinent shortening of the effective length, $L_g^{(eff)}$ of the gated portion of the HEMT channel [3]. This effect leads to an additional sensitivity of the instability to the source-drain voltage.

The self-excitation of the plasma oscillations in HEMTs can occur when $\mathrm{Im}\,\omega > 0$, i.e., when u_g or/and u_d exceed some threshold values. If u_g and u_d are under the threshold ($\mathrm{Im}\,\omega < 0$), the DS and TT mechanisms can reveal in a decrease of $|\mathrm{Im}\,\omega|$ with u_g and u_d approaching their threshold values. This corresponds to a decrease in the plasma oscillations damping and sharpening the plasma resonances. Owing to this, in the case $u_g \ll u_d$, the recently observed resonant detection of THz radiation at room temperatures [20] might be attributed to the TT mechanism.

As shown a long time ago, the features of the electron high speed transport accompanied by the optical phonon emission [21]–[23] can lead to an additional mechanism of negative dynamic conductivity. If the voltage drop across the high-field region somewhat exceeds the intervalley separation, the mechanism of negative differential conductivity proposed in Ref. [24] can also contribute to the plasma instability in HEMTs. The inclusion of these effects might complicate the pattern of the plasma instability in HEMTs. This, however, requires further studies.

Our present treatment is based on the hydrodynamic electron transport model. The hydrodynamic description of the electron system in HEMTs with really high electron mobility is fairly adequate. This is because the electron-electron collisions dominate over the collisions of electrons on remote impurities. More general and precise modeling of the phenomena under consideration can be realized using ensemble Monte Carlo particle technique. However, the inclusion of rather strong electron-electron interactions leads to a significant complications in numerical calculations.

6. Conclusions

In conclusion, we have demonstrated that the electron TT effect in the saturation regime of HEMT operation might lead to the plasma instability at sufficiently short electron transit times in the high-field gate-drain region if the plasma resonant frequency and the TT frequencies are close to each other. The plasma instability associated with this mechanism can occur even at relatively low electron drift velocities in the gated section of the HEMT channel when the DS mechanism is weak. The self-excitation of the plasma oscillations in HEMTs due to the combination of DS and TT mechanisms can used for the generation of THz radiation.

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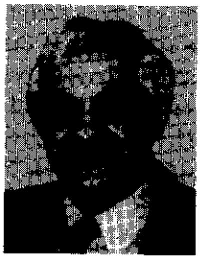

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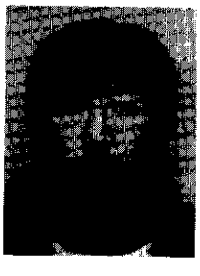

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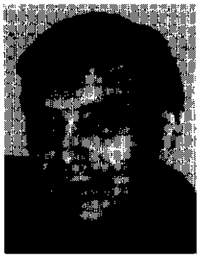

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