Transit-time resonances enabling amplification and generation of terahertz radiation in periodic graphene p-i-n structures with the Zener-Klein interband tunneling

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The Zener-Klein (ZK) interband tunneling in graphene layers (GLs) with the lateral n-i-n and p-i-n junctions results in the nonlinear I-V characteristics that can be used for the rectification and detection of the terahertz (THz) signals. The transit time delay of the tunneling electrons and holes in the depletion regions leads to the phase shift between the THz current and THz voltage causing the negative dynamic conductance in a certain frequency range and resulting in the so-called transit-time (TT) instability. The combination of the ZK tunneling and the TT negative dynamic conductance enables resonant THz detection and the amplification and generation of THz radiation. We propose and evaluate the THz devices based on periodic cascade GL p-i-n structures exhibiting the TT resonances (GPIN-TTDs). Such structures can serve as THz amplifiers and, being placed in a Fabri-Perot cavity, or coupled to a THz antenna or using a ring oscillator connection, as THz radiation sources.

I. INTRODUCTION

Since the prediction of the Dyakonov-Shur plasma instability in the current-driven two-dimensional electron channels of the field-effect transistor structures leading to the terahertz (THz) emission [1, 2], such an emission have been extensively discussed in the literature (see, for example, [3–15] and the references therein). As demonstrated recently [16–18], the Coulomb electron drag [19–23] by the injected hot electrons and holes in the graphene layer (GL) n-i-n and p-i-n structures [1, 2, 15, 24, 25] (see also [26]), could enable a specific instability mechanism. In these structures, the gated regions play the role of the plasmonic resonant cavities. The finite transit time of the holes and electrons injected or generated in the depleted regions of the GL channel can enable the negative dynamic conductivity on the order of the inverse transit time [27, 28].

In this paper, we propose and analyze the periodic ungated GL p-i-n transit time (TT) device (GPIN-TTD) with the i-region exhibiting the negative dynamic conductivity due to the TT effect. The generation of the holes and electrons in the i-region is associated with the Zener-Klein (ZK) interband tunneling [29–32]. In contrast to our previous works on the TT effects in GL structures, here we assume the following:

- The ungated d-regions are chemically doped. The plasma (PL) frequency of these regions is markedly larger than that of the gated regions and the ungated PL resonances are out of the frequency range under consideration;
- The Coulomb carrier drag in the relatively long d-regions in the GPIN-TTDs studied in [13, 23] is weak due to a short momentum relaxation time of the carriers injected into these regions from the i-region, especially at elevated voltage biases when the emission of optical phonons is essential.
- The combination of the ZK interband tunneling and the TT effects, which, in particular, enables the rectification and detection of the impinging THz radiation [33], might provide the THz radiation amplification and generation.

As we show below, the TT effects in the GPIN-TTD can result in the negative real part of its impedance enabling the amplification of the THz radiation. In the GPIN-TTDs with the optimized parameters, the TT resonances might be associated with the interplay of the i-region capacitance and the d-regions inductance. This can lead to the self-excitation of the THz oscillation and generation of the THz radiation in the GPIN-TTDs supplied by a proper antenna or incorporated into a ring oscillator configuration. Since the overall dimension of the periodic structure could be comparable to the THz radiation wavelength, it could also take a role of a distributive coupling antenna structure. Alternatively, the GPIN-TTDs can serve as an active media in the THz sources using Fabri-Perot cavities.
II. DEVICE MODEL

We consider the periodic GPIN-TTD structure based on a GL channel embedded in a dielectric, for example, hBN or polyimide, or placed on a dielectric substrate and having a thin top passivating layer. The GPIN-TTD structure can also be suspended in air or vacuum. The GL channel comprises a periodic array of the undoped i-regions of the length $2l$ placed between the chemically doped p- and n-regions (d-regions) of the lengths $l^p$ and $l^n$ ($l^p = l^n = L$). The channel is bounded by the side source and drain contacts. The drain-to-source bias voltage $\overline{V}$ results in alternating lateral reverse biased p-i-n junctions (with the depleted i-regions), p- and n-quasi-neutral region, and the forward biased n-i-p junctions (with the i-regions in these junctions, which below are referred to as the f-regions, filled with the injected electrons and holes and having a high conductance).

The bias voltage, $\overline{V}$, per one period (section) of the structure (assuming that the sections are equivalent) is equal to $\overline{V} = \overline{\Phi}/N$, where $N$ is the number of sections. Figure 1 shows the GPIN-TTD structure (the top passivating layer is not shown) and the potential profile of the p-i-n junction in one of the reverse-biased device sections. The GPIN-TTD can be connected with a THz antenna, can serve as a distributed antenna, form a ring oscillator structure, or can be placed in the Fabri-Perot cavity.

We assume that the holes and electrons generated in the i-regions due to the ZK interband tunneling propagate ballistically or near ballistically [34–40]. Due to the specific of the ZK interband tunneling in GLs [29], the velocity, $v = \pm v_W$, of the generated electrons and holes is directed primarily along the electric field in the i-region, i.e., in the $x-$ direction. If the i-region length $2l \lesssim 1 \mu$m, the transport of the injected holes and electrons in the i-region can be ballistic at room and lower temperatures [40, 41]. In this case, the average electron and hole velocities across the entire i-region are equal to $\pm v_W$, where $v_W \approx 10^8$ cm/s is the characteristic carrier velocity in GLs.

If the voltage drop across the i-regions between the d-regions $\Phi$ includes the AC component $\delta \Phi_\omega$, i.e., $\Phi = \overline{\Phi} + \delta \Phi_\omega \exp(-i\omega t)$ (where $\delta \Phi_\omega$ and $\omega$ are the THz signal amplitude and frequency), the holes and electrons generated due to tunneling and propagated in the i-regions (and the GL channel as a whole) induce the terminal AC current. Due to the finite time of the electron transit, the induced AC current and the AC voltage could have the opposite phase. The latter implies that the dynamic conductance of the i-region can be negative.

The impedance of the periodic GPIN-TTD structure $Z_{\omega}^{GPIN-TTD}$ is the sum of the impedances of the sections $Z_{\omega}^i$, i.e., $Z_{\omega}^{GPIN-TTD} = NZ_{\omega}$ with

$$Z_{\omega} = Z_p^i + Z_n^i + Z_w^i + Z_f^i,$$

where $Z_p^i$, $Z_n^i$, $Z_w^i$, and $Z_f^i$ are the impedances of the pertinent regions. For $N$ equivalent sections $Z_{\omega}^{GPIN-TTD} = NZ_{\omega}$.

Below we calculate the frequency-dependent impedance of the GPIN-TTD as a function of the structural characteristics. The numerical calculations were performed using Maple 2021 software (Maplesoft, Waterloo, ON, Canada).

III. THE IMPEDANCE OF THE GPIN-TTD

For the most interesting voltage range $e\overline{\Phi}_i > T$, where $T$ is the temperature in the energy units, we obtain for the AC current density (the current per unit width of the device in the direction along the gate edges)

$$\delta J_\omega = \sigma_\omega \delta \Phi_\omega. \quad (2)$$

For for the blade-like p- and n-regions and the electric field concentrated near these regions [27], the AC conductance accounting for the frequency-dependent factor associated with the TT effect (with the finiteness of the hole and electron TT, $t^i = 2l/v_W$, across the i-region) is given by

$$\sigma_\omega^i = \sigma_0 \exp(i\omega t^i/2) - i\omega e^i. \quad (3)$$
The differential DC conductance of the i-region \( \sigma^i \) determined by the DC voltage drop, \( \Phi \), across this region, and the geometrical capacitance of this region are

\[
\sigma^i = \frac{b e^2}{h} \sqrt{\frac{e \Phi}{2 \hbar \nu W}}
\]

and

\[
\epsilon^i = \frac{\kappa_{eff} \Lambda}{2 \pi^2} \propto \kappa_{eff}, \tag{5}
\]

respectively. In Eqs. (4) and (5), \( e = |e| \) is the carrier charge, \( b = [3 \Gamma(1/4) \Gamma(1/2)/2 \Gamma(3/4)] \approx 0.0716 \) \cite{28}, \( \kappa_{eff} \) is the effective dielectric constant accounting for the dielectric constant, \( \kappa_S \) and \( \kappa_T \), of the substrate and top/passivation layer at the frequency in the THz range and their thickness, \( \Lambda = (L/l) \tan^{-1} \left( \frac{1}{\sqrt{(L/l)^2 - 1}} + \ln\left[(L/l) + \sqrt{(L/l)^2 - 1}\right] \right) \) \cite{41} (see also \cite{42–45}), is a logarithmic factor depending on the device geometrical parameters (capacitance factor), and \( J_0(s) \) is the Bessel functions. Generally, \( \kappa_{eff} = (\kappa_S + \kappa_T)/2 \). In the case of a rather thin top/passivation layer with the thickness \( w \ll 2l \), \( \kappa_{eff} \approx (\kappa_S + 1)/2 \). The factor of 1/2 is because the streamlines of electric field emerge out of the top layer.

Examples of \( \Lambda \) and \( \epsilon^i \) as functions of the i-region length \( 2l \) are shown in Fig. 2.

The Drude formula yields the following expression for the AC conductance of the d-region:

\[
\sigma^d = \frac{i \nu}{\omega + i \nu} \sigma^d, \tag{6}
\]

where

\[
\sigma^d = \frac{e^2 \mu}{\pi \hbar^2 L \nu} \propto \frac{\mu}{L \nu} \tag{7}
\]

\begin{figure}[h]
\centering
\includegraphics[width=0.8\textwidth]{fig2.png}
\caption{Capacitance factor \( \Lambda \) versus i-regions length \( 2l \) for different d-regions length \( L \) and i-region capacitance \( \epsilon^i \) (inset) for \( L = 1.0 \, \mu m \) and different values of dielectric constant \( \kappa_{eff} = 1, 1.5, 2, \) and 2.5.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=0.8\textwidth]{fig3.png}
\caption{Net DC resistance of highly conducting regions \( 2I/\sigma_d \) versus normalized Fermi energy \( \mu/T \) for different values of parameter \( \eta = \nu L/\omega W \) (\( T = 300 \, \text{K} \)). Inset shows \( \pi/T \) as a function of parameter \( \eta \).}
\end{figure}

is the doped region DC conductance, \( \nu \) is the collision frequency of the carriers in the d-regions with phonons, impurities (remote and residual), and defects, and \( \mu \) is the carrier Fermi energy in these regions. Equation (7) accounts for the kinetic inductance associated with the quasi-equilibrium holes and electrons in the d-region.

Considering Eqs. (6) and (7) and accounting for the d-regions capacitance, we obtain the d-region impedance:

\[
Z^d \approx \frac{1}{\sigma^d \nu^2} = \frac{1}{\sigma^d \nu^2} + \frac{\eta L}{\omega c^d \Omega_{PL}}, \tag{8}
\]

where \( \Omega_{PL} = 1/\sqrt{c^d \nu^2} \) and \( c^d \) is the d-region inductance and \( \eta L/\omega c^d \Omega_{PL} \) is the d-region impedance. For \( \nu < \omega \ll \Omega_{PL} \), the frequency \( \Omega_{PL} \) can be fairly large compared to the gated graphene structures with the same doping and length due to a relatively small capacitance of the ungated d-regions.

In the following, we assume that the signal frequency \( \nu < \omega \ll \Omega_{PL} \). In this case, accounting for the f-region impedance, for the net impedance of the highly conducting regions (the d- and f-regions) in the periodic structure section we obtain: (see Appendix A)
\[ 2Z_d^f + Z_\omega \simeq \frac{2}{\sigma^4} F - i \frac{2}{\sigma^4} \frac{\omega}{\nu}. \] \tag{10}

Here the factor
\[ F = 1 + \left( \frac{\pi v_W}{8L\nu} \right) \frac{\mu}{(\mu \ln 2 + \pi^2 T/12)} \exp \left( \frac{\mu}{T} \right) \] \tag{11}

reflects the contribution of the f-regions to the GL channel resistance. In Eq. (10), we disregarded the reactive component of the f-region conductance.

Accounting for Eqs. (1) and (10), we obtain the net impedance, \( Z_\omega \), of one of the periodic structure sections
\[ Z_\omega = \frac{2}{\sigma^4} F - i \frac{2}{\sigma^4} \frac{\omega}{\nu} + \frac{1}{\pi^2 e^{\pi^2/2} J_0(\pi L/2) - i \omega e^{\pi L/2}}. \tag{12} \]

In the low-frequency limit, we obtain from Eq. (12) \( Z_0 = 1/\sigma^4 + 2F/\sigma^4 \).

The quantity \( 2F/\sigma^4 \) exhibits a minimum as a function of \( \mu \) at \( \mu = \bar{\mu} \sim T \) with \( \bar{\mu} \) determined by the parameter \( \eta = (\nu L/v_W) \). For \( L = (0.5 - 1.5) \mu m \) and \( \nu = 1 \) ps, one obtains \( \eta \simeq 0.5 - 1.5 \). Figure 3 shows the dependence of the net DC resistance of the highly conducting parts of the channel (comprising the d- and f-regions) on the normalized Fermi energy \( \mu/T \) calculated for different values of the parameter \( \eta \). The inset shows the Fermi energy, \( \bar{\mu} \), corresponding to the net resistance minimum, as a function of \( \nu \).

In the following, we chose \( \mu = \bar{\mu} \). Table I provides \( \min (2F/\sigma^4) \) and \( \min (2/\sigma^4) \) calculated for several values of the parameter \( \eta \). One can see from Table I that the resistance of the f-regions is crucial because \( \min (2F/\sigma^4) > \min (2/\sigma^4) \).

For the real and imaginary parts of the GPIN-TTD section Eq. (12) yields:

\[ \text{Re} Z_\omega \simeq \frac{2}{\sigma^4} F + \frac{G(\omega t^1/2)}{\sigma^4 [G^2(\omega t^1/2) + H^2(\omega t^1/2)]}, \tag{13} \]
\[ \text{Im} Z_\omega \simeq -\frac{2}{\sigma^4} \frac{\omega}{\nu} - \frac{H(\omega t^1/2)}{\sigma^4 [G^2(\omega t^1/2) + H^2(\omega t^1/2)]}. \tag{14} \]

Here
\[ G(\pi X) = J_0(\pi X) \cos(\pi X), \tag{15} \]
\[ H(\pi X) = J_0(\pi X) \sin(\pi X) - \pi X \xi^i, \tag{16} \]
where \( X = \omega t^1/2 \pi \) and \( \xi^i = (2e^i/t^i \sigma^4) \). The first terms on the right-side of Eqs. (13) and (14) reflect the resistance and kinetic induction of the highly conducting regions, respectively.

IV. SPECTRAL CHARACTERISTICS OF THE IMPEDANCE: QUALITATIVE ANALYSIS

Figure 4(a) shows the functions \( G(\pi X) \) and \( H(\pi X) \) (inset). As seen, \( G(\pi X) = G(\omega t^1/2) \) is negative in certain frequency ranges \( \omega < \Omega_m^0 \), with \( \Omega_m^0 \) and \( \Omega_m^1 \)
obeying the equations $G(\Omega_m^0 t^i/2) = 0$ and $G(\Omega_m^1 t^i/2) = 0$, respectively. The function $G(\pi X)$ exhibits minima at $X = X_m$, where $m = 1, 2, 3, \ldots$ and $X_1 \approx 0.632$, $X_2 \approx 1.632, X_3 \approx 2.632, \ldots$. In these minima $G(\pi X_m) = G_m < 0$. The minima correspond to the characteristic TT frequencies $\Omega_m^N/2\pi = X_m/t^i = X_m v_W/2l$, where $\Omega_m^0 < \Omega_m^N < \Omega_m$. Figure 4(b) shows the frequencies $\Omega_m^N/2\pi$ as functions of the i-region length $2l$. For sufficiently short i-regions ($2l \lesssim 1 \mu m$), $\Omega_m^0/2\pi, \Omega_m^1/2\pi$, and $\Omega_m^N/2\pi$ are close to zero, and the contribution $\Omega_m^N/2\pi$ to the THz range.

Since in certain frequency ranges $G(\pi X) = G(\omega t^i/2) < 0$, the impedance real part $\text{Re} \ Z_\omega$ given by Eq. (13) can be negative provided that the net d- and t-region resistance $2F/\sigma^d$ is sufficiently small.

The function $H(\pi X)$ can change the sign if $d[J_0(\pi X) \sin(\pi X)]/dX|_{X=0} > \pi \xi^d$. This is possible if $\xi^d < 1$. If $H(\omega t^i/2)$ is close to zero, and the contribution of d-regions inductance, reflected by the first term on the right-hand side of Eq. (14), is small, the impedance imaginary part $\text{Im} \ Z_\omega$ can turn to zero. This can occur at $\omega = \Omega_m^R > \Omega_m^0$. The frequencies $\Omega_m^R$ can be referred to as the TT resonant frequencies. Depending on the structure parameters, $\Omega_m^R$ can be larger than $\Omega_m^0$, but smaller than $\Omega_m^1$. In this case, $\text{Im} \ Z_\omega$ turns to zero in the frequency range in which $\text{Re} \ Z_\omega < 0$. In this case, the self-excitation of the THz oscillations in the GPIN-TTD might be possible [46]. Thus, the GPIN-TTD impedance frequency dependence is different whether $\Omega_m^R$ falls into the interval $(\Omega_m^0, \Omega_m^1)$ or not (with $\Omega_m^0 < \Omega_m^N < \Omega_m$).

For example, setting $2l = 0.3 - 0.5 \mu m, L = 0.5 \mu m, \mu = \overline{\mu} = 1.22T$, and $\overline{\Phi} = 100$ mV, we obtain $\sigma^i = (3.6 - 2.8) \text{ps}^{-1}$, $\min(2F/\sigma^d) = 0.069 \text{ps}$, $\Lambda = 2.89 - 2.36$, and $\omega/2\pi = \Omega_m^1/2\pi \approx 2.11 - 1.26 \text{THz}$ with $G_m \approx -0.1$ [see Fig. 4(a)]. In this case, the inequality $\xi^d < 1$ is satisfied for $k_{eff} < 1.18 - 2.125$. At $\overline{\Phi} = 400$ mV, instead of the latter inequality one obtains $k_{eff} < 2.36 - 4.25$. For longer i- and d-regions ($2l = 1.0 \mu m, L = 1.0 \mu m$, when $\sigma^i = (2.0 - 4.0) \text{ps}^{-1}$, $\min(2F/\sigma^d) = 0.0927 \text{ps}$, and $\Omega_m^1/2\pi \approx 0.632 \text{THz}$), the condition in question is more liberal: $k < 3.09 - 6.18$. However, since the ratio $|G_m|/\Omega_m^2$ steeply drops with increasing index $m$, the above condition can be realized for $m = 1$ only.

As follows from Eq. (14), at $\omega = \Omega_{TT}$, where

$$\Omega_{TT} = \sqrt{\frac{\sigma^d \nu}{2e^i}} = \frac{1}{\sqrt{L^d e^i}}. \tag{17}$$

the impedance imaginary part turns to zero.

Comparing the characteristic frequencies $\Omega_{PL}, \Omega_{TT}$, and $\Omega_m^N$, for their ratio we obtain

$$\frac{\Omega_{PL}}{\Omega_{TT}} = \sqrt{\frac{8\pi^2 e^i}{\kappa}} = \sqrt{\frac{4\ln\Lambda}{3}}. \tag{18}$$

$$\frac{\Omega_m^N}{\Omega_{TT}} = \frac{X_1 h v_W}{e(2l)} \sqrt{\frac{8\pi^3 L e^i}{\mu}} = \frac{X_1 h v_W}{e(2l)} \sqrt{\frac{4\pi L e^i}{\mu}}. \tag{19}$$
Setting as in the above estimates, $2l = (0.3 - 0.5) \mu m$, $L = 0.5 \mu m$ ($L = 2.89 - 2.364$), $\kappa = 1.5 - 2.5$, and $\mu = \pi = 1.22T$, we obtain $\Omega_{PL}/\Omega_{TT} \simeq 3.51 - 3.03$ and $\Omega_{N}^{\pi}/2\pi = 2.11 - 1.26$ THz with $\Omega_{N}^{\pi}/\Omega_{TT} \simeq 1.01 - 0.55$ (for $\kappa_{eff} = 1.5$) and $\Omega_{N}^{\pi}/\Omega_{TT} \simeq 1.3 - 0.71$ (for $\kappa_{eff} = 2.5$).

The estimates demonstrate that at realistic device parameters, Re $Z_\omega$ can be negative in the THz frequency range. Moreover, in some cases, at the frequencies corresponding to Re $Z_\omega < 0$, the impedance imaginary part Im $Z_\omega$ can turn to zero. As seen from these estimates, the disregarding of the PL response assumed above is justified up to the frequencies $\omega \sim \Omega_{N}^{\pi}$ (at least for not too short i-region when $\Omega_{1} \ll \Omega_{PL}$). The ratio $\Omega_{1}/\Omega_{PL}$ is small because the ungated d-regions capacitance is smaller than the i-region capacitance.

V. SPECTRAL CHARACTERISTICS OF THE IMPEDANCE: NUMERICAL CALCULATIONS

Figures 5 - 7 show the dependences of the GPIN-TTD impedance real and imaginary parts, Re $Z_\omega$ and Im $Z_\omega$, on the signal frequency $f = \omega/2\pi$ calculated numerically using Eq. (12) [or Eqs. (13) - (16)] for different parameters $2l$, $L$ (or $\eta = \nu L/v\omega$), and $\kappa_{eff}$, and for different values of DC voltage drop, $\Phi$, across the i-region. It is assumed that $T = 300$ K and $\nu = 1$ ps. Other parameters are taken from Table I.

As seen from these figures, Re $Z_{\omega}$ of the GPIN-TTDs with the chosen parameters is negative in certain frequency ranges around the frequency $\Omega_{N}^{\pi}/2\pi$, i.e., the frequencies corresponding to the TT fundamental resonances. At some parameters and elevated voltage $\Phi$, Re $Z_{\omega} < 0$ also near the second resonance with the frequency about $\Omega_{N}^{\pi}$ [see Figs. 5(b) and 6(b)]. The negativity of the impedance real part can be accompanied by the zero imaginary part. As follows from Figs. 5 - 7, this occurs at the frequency close to the fundamental TT resonance, but not in the case at the second resonance [see Figs. 5(b) and 6(b)]. To illustrate this, Figs. 8(a) and 8(b) show the plots taken from Figs. 5(a), 6(a) and 5(b), 6(b) corresponding to $\kappa_{eff} = 2$ and the frequencies $\Omega_{N}^{\pi}/2\pi = 0.63$ THz and $\Omega_{N}^{\pi}/2\pi = 0.84$ THz, respectively.

VI. ABSORPTION/AMPLIFICATION AND GENERATION OF INCIDENT THZ RADIATION

When the i-region Re $\sigma_{i}^{\nu} > 0$, so that Re $Z_{\omega}$ is positive, the GPIN-TTD absorbs the incident radiation. If Re $\sigma_{i}^{\nu} < 0$ in a certain frequency range, Re $Z_{\omega}$ can be negative in this range. In the latter case, the GPIN-TTD can serve as an active region of the devices generating electromagnetic radiation, particularly, in the THz range. Thus, the GPIN-TTD exhibiting Re $Z_{\omega} < 0$ and supplied with an antenna can be the THz radiation source. The pertinent conditions are Im $Z_{\omega} = 0$.
and $\text{Re } Z_{\omega} H + R^A = 0$, where $R^A$ is the antenna radiation resistance (see, for example, [40]) and $H$ is the GPIN-TTD width. Assuming that in line with Figs. 5 - 8, at the frequency corresponding to $\text{Im } Z_{\omega} = 0$, $\text{Re } Z_{\omega} \simeq -2.5 \text{ ps (} \sim 2.25 \Omega \cdot \text{cm})$, $H = 10^3 \mu \text{m}$, for the periodic GPIN-TTDs with $N = 10$ sections (periods), one obtains $\text{Re } N Z_{\omega}/H \simeq -225 \Omega$. This implies that the impedance matching requires an antenna with sufficiently large radiation resistance like those considered in [47, 48].

Sufficiently long GPIN-TTD without a special antenna (having only the highly conducting side contacts) with the impedance matched to the free space impedance $Z = 4\pi/e \simeq (4\pi/3)10^{-10} \text{ s/cm}$ (in the Gaussian units or $Z = 120\pi \Omega \simeq 377 \Omega$), can effectively emit the THz radiation. For the parameters as in the above estimate, this requires the number of the structure sections $N \simeq 16 - 17$.

The GPIN-TTD with $\text{Re } Z_{\omega} < 0$ can amplify the incident THz radiation (the absorption coefficient is negative). In this situation, the THz source can consist of such a GPIN-TTD placed into a Fabri-Perot cavity. The coefficient, $A_{\omega}$, of the normally incident THz radiation absorption in the GPIN-TTD channel is expressed via the impedance $\tilde{Z}_{\omega}$ as (see, for example, [49, 51])

$$A_{\omega} = 1 - \left| \frac{\tilde{Z}_{\omega}/Z - 1}{\tilde{Z}_{\omega}/Z + 1} \right|^2 \left( \frac{4\text{Re } \tilde{Z}_{\omega}/Z}{(\text{Re } \tilde{Z}_{\omega}/Z + 1)^2 + (\text{Im } \tilde{Z}_{\omega}/Z)^2} \right),$$  \hspace{1cm} (20)

where $\tilde{Z}_{\omega} = Z_{\omega}\sqrt{\kappa_{eff}}/D$ and $D = 4l + 2L$ is the section length.

First of all, we estimate $A_{\omega}$ for the frequency corresponding to $\text{Im } Z_{\omega} = 0$, setting as in the above estimate $\text{Re } Z_{\omega} \simeq -(2.5 - 5.0) \text{ ps}$, $D = (2.0 - 4.0) \mu \text{m}$, $\kappa_{eff} = 1 - 3$, and accounting that $|\tilde{Z}_{\omega}/Z| \gg 1$, we obtain from Eq. (20)

$$A_{\omega} \simeq \frac{4Z}{\text{Re } \tilde{Z}_{\omega}} \simeq -(5 - 23) \%. \hspace{1cm} (21)$$

Figure 9 shows the absorption coefficient $A_{\omega}$ versus the frequency of the incident THz radiation $\omega/2\pi$ calculated using Eqs. (12) and (20) for the GPIN-TTD parameters corresponding to the plots in Fig. 8. As seen, the absorption coefficient $A_{\omega}$ exhibits a resonant plasmonic peak at the frequencies close to the TT characteristic frequency $\Omega_1^N$, i.e., to the fundamental TT resonance. The estimate given by Eq. (21) is in line with the results of numerical calculations shown in Fig. 9.
FIG. 8. The frequency dependences of the GPIN-TTD impedance real (solid lines) and imaginary parts (dashed lines) for \( \kappa_{eff} = 1, 2, \) and 3 and other structural parameters corresponding to the TT characteristic frequencies \( \Omega_N^{\gamma}/2\pi = 0.63 \) THz and \( \Omega_N^{\gamma}/2\pi = 0.84 \) THz (shown also in Figs. 5 and 6): (a) \( \Phi = 100 \) mV and (b) \( \Phi = 400 \) mV.

VII. COMMENTS

(i) At elevated bias voltages, the potential drop across the i-region can be so large that \( e\Phi > \hbar \omega_0 \), where \( \hbar \omega_0 \approx 200 \) meV is the optical phonon energy in GLs. In this situation, a fraction of the electrons and holes generated in the i-region spontaneously emits the optical phonons. This leads to a loss of these carriers energy and momentum. As a result, the directed velocity of the carrier emitted the optical phonons and the average directed velocity \( <v_x> \) become smaller than \( \pm v_{TP} \). Hence, the optical phonon emission at elevated voltages can result in an increase of the carrier TT and a decrease in the characteristic TT frequency that leads to a red shift of the impedance and absorption coefficient peaks, although a rather moderate [about of (3 – 5)%] for the parameters corresponding to the plots in Figs. 6(b) and 7(b) (see Appendix B).

(ii) The realization of the negative impedance in the GPIN-TTD and the amplification of the THz radiation passing its plane, associated with the TT resonances, require low or moderate values of the THz dielectric constants of the substrate and the top/passivating layer. Among such materials one can mention crystalline hBN as a material for the substrate with a thin passivation layer \( \kappa_{eff} \approx 3 \) [52] as well as polyimide, porous BN/polyimide composite, and amorphous HBN \( \kappa \approx 1.16 \) [53, 54]. It is interesting that the GPIN-TTDs on some of such substrates can be used for flexible THz sources (see [53]).

(iii) When the DC voltage drop across the i-region \( \Phi \) is relatively large, the space charge of the tunnel-generated carriers [disregarded above, in particular, in Eq. (4)] can play some role modifying the DC potential spatial distribution [28]. This, in turn, leads to a variation (to some increase as shown by numerical simulations [28]) of the tunnel current. The space-charge effects in the structures under consideration are characterized by the parameter \( \gamma_{SC} = [e^2/4\pi^2 \kappa_{eff} \hbar^{3/2} v_N^{3/2} W]/2e\Phi \) [28]. The effects in question become crucial when \( \gamma_{SC} \geq 1 \). The parameter \( \gamma_{SC} \) reflects the ratio of the electric field formed by the hole and electron charges in the i-regions and the electric field created by the DC voltage between the p- and n-regions. Setting \( 2l = 0.5 – 0.75 \) \( \mu m \) and \( \kappa_{eff} = 2 – 3 \) for \( \Phi = (100 – 400) \) mV, we obtain \( \gamma_{SC} \approx 0.055 – 0.101 \). This implies that in the GPIN-TTDs at elevated voltages \( \Phi \gtrsim 400 \) mV), the space-charge effects can be pronounced. However, such effects increase the differential conductance \( \sigma' \) and, therefore, increase of the absolute value of the negative dynamic conductance. The latter is useful for reinforcing the THz radiation amplification and generation. The GPIN-TTD structures with multiple non-Bernal (twisted) GLs (instead of a single GL in the devices considered above) also can be used as the THz sources. In these multiple-GPIN-TTDs the negative conductances of the i-regions of each GL add up,
while the net geometrical capacitance is close to that for single-GL structures. Due to the net space-charge in the parallel i-regions, their net conductance can be rather marked affecting the device performance. In general, the multiple-GPIN-TTDs can surpass the single-GPIN-TTDs, although their analysis requires a separate study.

(iv) It is interesting to compare the amplification coefficient estimate given by Eq. (21) with the interband amplification coefficient of the GL with the optical or injection pumping exhibiting the interband population inversion. The amplification coefficient maximum of the latter is equal \( A_{\text{int}}^{\text{max}} \approx \pi e^2 / \hbar c \sqrt{\kappa_{\text{eff}}} \). This yields max \( (A_{\omega}^{\text{int}}) \approx (1.3 - 2.3)\% \), i.e., markedly smaller than the GPIN-TTD maximum amplification coefficient. In contrast to the pumped GLs, which can amplify the THz radiation in a fairly wide frequency range, the GPIN-TTDs exhibit the amplification at frequencies close to the TT resonances. The GPIN-TTD structures comprise the reverse- and forward-biased p-i-n junctions. Similar situation occurs in the dual-gated structures. The p-i-n junctions of both types can exhibit the THz radiation amplification but associated with different physical mechanisms. The co-existence of these mechanisms can reinforce the amplification effect.

(v) The values of the amplification coefficient, \( -A_{\omega} \), demonstrated above correspond to nonoptimal match of the GPIN-TTD impedance and the free space impedance. The device parameters used in the above calculation correspond to large values of \( |\text{Re} \tilde{Z}_\omega / Z| \propto |\text{Re} Z_\omega| \). Properly choosing the conductance of the d- and f-regions (by varying the doping level and the d-region length), one can, in principle, achieve the condition \( \text{Re} \tilde{Z}_\omega / Z \sim 1 \) (with \( \text{Re} Z_\omega < 0 \)). One can expect that such an optimization can provide a fairly large amplification coefficient and even to achieve \( A_{\omega} > 1 \).

(vi) Above, in particular, in Eq. (1), we neglected the resistance of the side contacts. This is justified by the possibility to form very low resistance contact to GLs, as low as 0.0045 Ω·cm [63] and 0.02 Ω·cm [65] for Au contacts. Hence, the contact resistance can be much smaller (by two or more orders of magnitude) than the impedance of the forward-biased n-i-p junctions \( Z_f \sim (4 - 5) \) Ω·cm.

(vii) When the PL and TT frequencies are comparable, the interaction of the pertinent resonances can add complexity to the pattern of the phenomena under consideration. In contrast to the situation analyzed above, this can happen primarily in the GPIN-TTDs with the gated p- and n-regions, which are beyond the scope of this paper.

VIII. CONCLUSIONS

We proposed GPIN-TTDs based on periodic ungated GL structures with the p-i-n junctions exhibiting the TT
effects associated with the holes and electrons generated in the reverse-biased i-regions due to the Zener-Klein interband tunneling. The proposed GPIN-TTDs can be used for the amplification and generation of the THz radiation. Their operation is associated with the TT resonances in certain ranges of frequencies resulting in the negative dynamic conductance of the i-regions. The TT resonant frequencies are determined by the i-region capacitance and the kinetic inductance of the tunnel-generated carriers in the i-regions as well as the carriers in the doped p- and n-regions. Since the GPIN-TTDs can demonstrate rather strong amplification of the THz signals, the THz sources based on the GPIN-TTDs coupled to an antenna (or being used as distributed coupling antenna structure), forming the ring circuit configuration, or placed within a Fabri-Perot cavity, can markedly surpass the emitters using the interband negative conductance caused by optical or injection pumping.

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APPENDIX A.

The impedance of forward-biased n-p junctions

The current via the undoped regions between the p- and n-regions forward-biased lateral p-i-n junctions in the GPIN-TTDs (we refer to this region as the f-regions) comprises the tunneling and thermionic components. The former is characterized by the differential conductance similar to that for the reversed-bias n-i-p junction given by Eq. (4), but with the voltage, $\Phi_f$, markedly smaller than $\Phi$. Due to this, the tunneling component is much smaller than the thermionic component. Considering that the electrons and holes injected from the n- and p-regions into the p- and n-regions, respectively, have kinetic energies exceeding $\mu - e\Phi_f$ and taking into account the angular spread of the electron and hole velocities, the density of the thermionic current via the forward-bias p-i-n junction in the GPIN-TTDs under consideration can be presented as

$$J_f \approx \frac{4eT^2}{h^2eW} \left( \frac{\mu}{T} \ln 2 + \frac{\pi^2}{12} \right) \exp \left( -\frac{\mu}{T} \right) \times \left[ \exp \left( \frac{e\Phi_f}{T} \right) - 1 \right]. \quad (A1)$$

Considering the smallness of the voltage drop across the forward-biased n-p junction $\Phi_f$, and disregarding the kinetic inductance of this junction, from Eq. (A1) we obtain the following formula for the junction impedance:

$$Z^f_\omega \approx Z^f_0 = \frac{\pi^2h^2eW}{4e^2(\mu \ln 2 + \pi^2T/12)} \exp(\mu/T). \quad (A2)$$

In particular, accounting for Eq. (A2), we obtain

$$\text{Re}(2Z^f_\omega + Z^f_0) \approx \frac{2}{\sigma^2} \left[ 1 + \left( \frac{\pi v_W}{8\mu v_W} \right) \left( \mu \exp(\mu/T) \right) \right]. \quad (A3)$$

APPENDIX B.

Role of the optical phonons emission

Assuming for simplicity that the optical phonon emission makes the carrier distribution semi-isotropic, the average directed carrier velocity after the emission can be set as $v_x \approx \pm v_{W} / 2$. Hence, for the carriers in the i-regions where their energy $E \approx h\omega_0$ the average velocity is equal to $\pm v_{W}$ and for the carriers in the region where $E > h\omega_0$ this velocity is equal to $\pm v_{W} / 2$. Considering that the characteristic time, $\tau_0$, of the spontaneous optical phonon emission is finite, a portion of the carriers having the energy $E > h\omega_0$ can continue to propagate with the velocity $\pm v_{W}$. Summarizing all this, we arrive at the following rough estimate for $v_x >$:

$$\frac{< v_x >}{v_{W}} \approx \frac{1}{2} \left( 1 + \frac{l_0}{l} \right) + \frac{v_W \tau_0}{4l} \left[ 1 - e^{-2(l-l_0)/v_W \tau_0} \right]. \quad (B1)$$

where

$$l_0 = \frac{e\Phi - h\omega_0}{e\Phi} \cdot \exp \left( \frac{e\Phi - h\omega_0}{e\Phi} \right). \quad (B2)$$

When $v_W \tau_0 > 2l$, at $e\Phi > h\omega_0$ from Eqs. (B1) and (B2) for the average velocity $< v_x > / v_{W}$ and, consequently, for the average transit time $l_0 = 2l / < v_x > > t_W / < v_x >$ one obtains

$$\frac{< v_x >}{v_{W}} \approx 1 - \left( \frac{l-l_0}{2l_{W}\tau_0} \right)^2, \quad t_0 \approx t \left[ 1 + \left( \frac{l-l_0}{2l_{W}\tau_0} \right)^2 \right]. \quad (B3)$$

Assuming that $\tau_0 \approx 1$ ps (see, for example, [66, 67]), for $2l = (0.5 - 0.75) \mu \text{m}$ and $\Phi = 400 \text{ mV}$ this corresponds to the plots in Figs. 6(b) and 7(b) and using Eqs. (B2) and (B3), we find $(t_0 - t) / t \approx (3.1 - 4.7) \%$.

A low density of states near the Dirac point leads to a decrease of the optical emission probability in comparison with $\tau_0^{-1}$. Anisotropy of the carrier scattering on the optical phonons results in somewhat larger values of the directed carrier velocity compared to $v_{W} / 2$. This implies that the above estimates provide a conservative value of the average velocity and, hence, an overrated TT value.
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