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Graphene electronics for terahertz electron-beam radiation

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Abstract

By virtue of their distinctive electronic properties (including linear energy dispersion, large velocity, and potentially ultra-high mobility even at room temperature), charge carriers in single-layer graphene are uniquely suited to radiation mechanisms that so far have been the primary domain of electron beams in vacuum-based systems. Here, we consider the use of sinusoidally corrugated graphene sheets for the generation of THz light based on a fundamentally new cyclotron-like radiation process, which does not require the application of any external magnetic field. Instead, periodic angular motion under bias is simply produced by the graphene mechanical corrugation, combined with its two-dimensional nature which ensures that the carrier trajectories perfectly conform to the corrugation. Numerical simulations indicate that technologically significant output power levels can correspondingly be obtained at geometrically tunable THz frequencies. This mechanism (as well as similar electron-beam radiation processes such as the Smith–Purcell and Cherenkov effects in periodic nanostructures) may open the way for a new family of THz optoelectronic devices based on graphene, including solid-state ‘free-electron’ lasers potentially capable of room-temperature operation.

(Some figures may appear in colour only in the online journal)

1. Introduction

Light emission in semiconductors typically occurs via electronic transitions between different energy bands in bulk samples, or between quantum-confined states in low-dimensional systems. These emission processes form the basis of several well established devices operating at visible and infrared wavelengths, such as LEDs, diode lasers, and quantum cascade lasers. In recent years, their domain has also been extended to longer and longer wavelengths across the far-infrared spectrum through the development of terahertz quantum cascade lasers, whose operation however is fundamentally limited to cryogenic temperatures [1]. An alternative class of radiation mechanisms in solid-state systems involves the periodic acceleration of charge carriers under static bias conditions. This approach is similarly well established at microwave frequencies, as in the case of electronic oscillators such as tunneling diodes, transit-time devices, and Gunn diodes. In recent years, several efforts have also been devoted to devise similar processes suitable for THz light emission, motivated by fundamental questions about ultrafast carrier dynamics in low-dimensional systems as well as the emerging need for practical solid-state THz sources [2]. However, the resulting mechanisms mostly require cryogenic temperatures, and generally produce extremely broad and weak radiation spectra. These limitations are fundamentally related to basic material properties of traditional semiconductors, such as limited carrier velocities, low mobilities at room temperature, and broad thermal distributions of drift velocities.

The recent emergence of graphene, with its unique electronic and mechanical properties, offers a new perspective and fundamental benefits. As clearly established by extensive
experimental and theoretical work [3–7], charge transport in these single-layer sheets of carbon is dominated by two π-like energy bands whose separation is zero at the corners of the hexagonal Brillouin zone (commonly named the Dirac points). In the vicinity of these points, both bands feature a linear dispersion, \( E = \pm \hbar v_F k \), similar to the case of ultra-relativistic particles described by the massless Dirac equation, albeit at smaller speeds (\( v_F = 10^8 \text{ cm s}^{-1} \)). Large densities of electrons and holes can be created in these bands in a highly controllable fashion via capacitive effects in a gated configuration. With the same setup, exceptional electronic transport properties have been demonstrated, including micrometer-scale ballistic transport and record large room-temperature mobilities above \( 1 \times 10^5 \text{ cm}^2 \text{ V}^{-1} \text{ s}^{-1} \) in suspended samples [8, 9] and in samples deposited on atomically smooth, highly inert hexagonal boron nitride (h-BN) films [10, 11]. Potential device applications to THz photonics have also been investigated recently [12–21], mostly based on interband electronic transitions and related to the ultrafast carrier relaxation/recombination dynamics near the Dirac points. In particular, THz amplification [16] and spontaneous emission [17] under optical pumping have already been reported. The use of tunable plasmonic excitations in electrostatically gated graphene is also being widely investigated for several infrared (including THz) device applications [13, 15, 19, 22].

In this work we show that, by virtue of their distinctive electronic properties (including linear energy dispersion, relatively large velocity, and ultra-high mobilities), charge carriers in graphene are also well suited to radiation mechanisms that so far have been the primary domain of high-energy electron beams in vacuum-based systems. Such mechanisms include cyclotron emission, the Smith–Purcell effect [23, 24] and more in general Cherenkov radiation in photonic crystals [25], all of which involve charges traveling at constant speed in the presence of a periodic spatial modulation (in either their trajectory or the surrounding dielectric environment). In particular, here we investigate the use of sinusoidally corrugated graphene for the generation of THz light based on a cyclotron-like radiation process. The required geometry, illustrated schematically in figure 1, could be obtained by depositing a sheet of graphene on a grating patterned on the surface of a suitable substrate (e.g., an oxidized silicon wafer, as commonly employed in graphene electronics). Because of the ultrasmall thickness (i.e., single-atomic-layer) of graphene, conformal adhesion to the corrugated surface can be expected even with sub-micron grating periodicities. In the presence of a dc voltage in the direction of the corrugation, electrons (or holes) introduced in the graphene sheet via field doping will therefore undergo periodic angular motion and correspondingly radiate.

This mechanism is well established in the context of vacuum electron-beam devices such as free-electron lasers [26]. However, it represents a novel paradigm for light emission in condensed matter, and the formal analogy between electrons and holes in graphene and relativistic charges provides a natural motivation for its study. The use of geometrical constraints (as opposed to the application of an external magnetic field) to obtain radiation via angular motion is also fundamentally new. Specifically, in the case of a free-electron laser sinusoidal carrier trajectories are obtained with a periodic array of magnets with alternating poles (the undulator). In the present context, the same periodic angular motion is instead produced by the graphene mechanical corrugation, combined with its two-dimensional nature that ensures that the carrier trajectories perfectly conform to the corrugation. Rigorous electromagnetic simulations based on the finite difference time domain (FDTD) method indicate that technologically significant output power levels can be obtained with this approach at geometrically tunable THz wavelengths, even at room temperature. The same geometry can also in principle be used to provide stimulated emission and optical gain (in analogy with the physics of free-electron lasers), and therefore can potentially enable the development of a promising new family of solid-state THz lasers.

2. Results and discussion

The basic radiation properties of charges in periodic angular motion can be investigated using classical electrodynamics [27]. Simple analytical expressions are obtained from the Larmor formula in the case of motion with constant speed \( v_0 \) in a sinusoidal trajectory of period \( \Lambda \) and amplitude \( A \), in the limit where \( v_0 \ll c \) and \( A \ll \Lambda \). Under these conditions, the radiation frequency \( f_{rad} \) and output power per charge-carrier \( P_{ch} \) can be approximated as follows:

\[
 f_{rad} = \frac{v_0}{\Lambda}, \tag{1}
\]

\[
 P_{ch} = \frac{4\pi^3 q^2 A^2 v_0^4}{3\varepsilon_0 c^3 \Lambda^4}, \tag{2}
\]

where \( q, \varepsilon_0, \) and \( c \) are the electron charge, permittivity of free space, and speed of light, respectively. Recently, a theoretical analysis based on these expressions has been applied to the case of charge transport in periodically corrugated SiGe heterojunctions [28]. Equations (1) and (2), while only approximate, can be used to substantiate the advantages of the unique electronic structure of graphene for this application. Due to the linear energy dispersion near the Dirac points, all conduction electrons and holes in graphene travel at the same speed \( dE/d(\hbar k) = v_F = 10^8 \text{ cm s}^{-1} \).
This value is substantially larger than the maximum drift velocities achievable in typical semiconductors (by about an order of magnitude), which is desirable due to the \( v_0^4 \) dependence of the emitted power (see equation (2)). At the same time, proportionally higher emission frequencies \( f_{\text{rad}} \) within the THz spectral region can be obtained with graphene using realistic corrugation periods \( \Lambda \) on the order of a few hundred nanometers (see equation (1)). Furthermore, in a traditional two-dimensional electron gas with a parabolic energy band, the electron speed varies (linearly) with wavenumber \( k \). As a result, in the geometry of figure 1, a much narrower distribution of the electron velocity component along the direction of the corrugation can be expected in graphene compared to conventional semiconductors, particularly under highly non-equilibrium conditions. Since the cyclotron emission frequency \( f_{\text{rad}} \) is proportional to this velocity component (\( v_0 \) in equation (1)), the end result is a narrower and therefore stronger emission spectrum.

For a more rigorous analysis, we begin by employing a commercial FDTD package (FDTD solutions by Lumerical [29]) to compute the radiation output of a single electron in a sinusoidally corrugated sheet of graphene. In order to apply the FDTD method to this problem, we take advantage of the formal equivalence between a time-varying distribution of electric dipoles (described by a polarization density \( \mathbf{P}(\mathbf{r}, t) \)) and the current density \( \mathbf{J}(\mathbf{r}, t) \). The Fourier transform of the equivalent dipole distribution \( \mathbf{P}(\mathbf{r}, t) \) is identical to that of the polarization density \( \mathbf{P}(\mathbf{r}, t) \) whose time derivative is equal to \( \mathbf{J}(\mathbf{r}, t) \). Once the equivalent dipole distribution described by this polarization density is identified, it can be readily modeled using the built-in dipole radiation sources of the FDTD simulation engine.

Specifically, we consider an electron in an arbitrary Bloch state of wavevector \( \mathbf{k} \) (measured from the nearest Dirac point), in the presence of a sinusoidal corrugation along the \( x \) direction having amplitude \( A \) and period \( \Lambda \) (see figure 1). In the plane of the graphene sheet, the electron travels at constant speed \( v_F = 10^6 \text{ cm s}^{-1} \) along the direction of \( \mathbf{k} \). Due to the corrugation, this direction varies as a function of position, while remaining at a fixed angle \( \theta \) with respect to the projection of the \( x \)-axis on the corrugated surface. In terms of the system of coordinates shown in figure 1, the electron trajectory can be written as

\[
\mathbf{r}_e(t) = \mathbf{x}_e(t) + \mathbf{y}(\mathbf{x}_e(t)) \tan \theta + \mathbf{z} A \sin (px_e(t)),
\]

where \( x_e(t) \) is the instantaneous position along the \( x \) direction, \( p = 2\pi/\Lambda \) is the corrugation wavenumber, and

\[
l(x) = \int_0^x \sqrt{1 + A^2 p^2 \cos^2(px)} \, \text{d}x
\]

is the arc length of the corrugation profile along the \( x \) direction. The variable \( x_e(t) \) can be calculated numerically as a function of time from the equation \( \dot{\mathbf{x}}_e(t) = v_F \), which simply expresses the condition that the electron speed is constant and equal to \( v_F \), regardless of the wavevector \( \mathbf{k} \) and the externally applied electric field. The current density carried by the electron is simply related to its trajectory as \( \mathbf{J}(\mathbf{r}, t) = -q \mathbf{e} \mathbf{r}_e(t) \delta (\mathbf{r} - \mathbf{r}_e(t)) \). To determine the equivalent polarization density, we Fourier transform \( \mathbf{J}(\mathbf{r}, t) \) and divide the result by \(-io\omega\), leading to the following expression:

\[
P(\mathbf{r}, \omega) = \frac{q}{iw} \left[ \hat{\mathbf{x}} + \hat{\mathbf{y}}'(x) \tan \theta + \hat{\mathbf{z}} A \cos(px) \right] \times \delta(y - l(x) \tan \theta) \delta(z - A \sin(px)) e^{i\omega t},
\]

where \( \hat{\mathbf{y}}'(x) = \text{d}l/dx \) and \( l(x) \) is the time instant when the electron position along the corrugation satisfies \( x_e(t) = x \).

Equation (5) describes a continuous distribution of electric dipoles linearly positioned on the charge trajectory, having position-dependent magnitude, phase, and direction of the dipole moment. In the FDTD simulations, this distribution is discretized into a collection of neighboring dipoles separated by a small distance \( \Delta x \) along the \( x \) direction, with \( \Delta x \ll \Lambda \). A schematic plot of such an ensemble, taken from the Lumerical graphical user interface, is shown in figure 2, where each dipole is represented by a circle centered about its position and by an arrow in the direction of its moment. Specifically, in each simulation we model a sinusoidal trajectory of finite length, containing exactly 60 periods of the corrugation with 40 equally spaced dipoles per period. A uniform surrounding dielectric medium (air) is assumed, which dramatically simplifies the required computational grid. The simulation domain has the shape of a cube centered on the dipole distribution with 600-\( \mu \text{m} \) side length, and perfectly-matched-layer boundary conditions are employed on all boundaries. All the simulation parameters were selected after extensive convergence tests.

In figure 3(a) we plot the calculated radiation spectra produced by the equivalent dipole distribution of equation (5), for different trajectories (i.e., different values of the angle \( \theta \)) on a sinusoidal corrugation of period \( \Lambda = 300 \) \( \text{nm} \) and amplitude \( A = 50 \) \( \text{nm} \). As explained in the preceding paragraphs, these results also describe light emission from a single electron traveling along the same trajectories at constant speed \( v_F \). In the case of motion along the direction of the corrugation (i.e., \( \theta = 0^\circ \)), the emission spectrum of figure 3(a) consists of a sharp peak centered at about 2.7 THz and a weaker feature at twice that frequency.
These emission lines can be simply interpreted as the result of constructive interference among the output fields of all the radiating dipoles in the equivalent dipole distribution (higher-order harmonics, not included in the figure, are also present). The peak emission frequency of 2.7 THz is smaller than the predicted radiation frequency \( f_{rad} = 3.33 \) THz based on the approximate formula of equation (1) for \( \Lambda = 300 \text{ nm} \) and \( \theta = 0^\circ \) (i.e., \( v_0 = v_F \)). The difference is due to the non-negligible corrugation amplitude of the trajectories considered in figure 3(a) (\( A = 50 \text{ nm} \)), whereas equation (1) assumes \( A \ll \Lambda \). When this assumption is not satisfied, the distance traveled per period is substantially larger than \( \Lambda \) and the oscillation time is proportionally longer than \( \Lambda/v_0 \). As the angle \( \theta \) is increased and the electron velocity component along the direction of the corrugation \( (v_0 \approx v_F \cos \theta) \) correspondingly decreases, the radiation spectrum is found to red-shift and decrease in peak value. These observations are consistent with equations (1) and (2), which predict a decrease in both emission frequency \( f_{rad} \) and output power \( P_{ch} \) with decreasing \( v_0 \).

The emission lines of figure 3(a) feature extremely narrow spectral widths, as small as a few ten GHz. In practice, substantial spectral broadening can be expected due to various electronic scattering processes, involving, e.g., charged surface states in the graphene substrate, phonons, and surface roughness, which are clearly not included in the FDTD simulations just described. Their impact can be quantified as a broadening of the emission lines by the amount \( \Delta f = v_F/\pi d \) (where \( d \) is the mean free path between consecutive collisions) \[30\], and a proportional lowering of their peak values. In figure 3(b) we show the radiation spectrum produced by an electronic ensemble in graphene in the presence of the 300-nm-period, 50-nm-amplitude corrugation considered so far, for different values of \( d \). These traces were obtained by adding up the single-electron spectra corresponding to different electron trajectories (such as the examples shown in figure 3(a)), each convolved with a Lorentzian lineshape function of full width at half maximum \( \Delta f \). More details of the computation method are presented below. It follows from this figure that pronounced emission spectra with reasonably large quality factors can be expected even with mean free paths of a few hundred nanometers. Such values can be achieved in high-quality graphene samples even at room temperature \[8, 11\], which is another key enabling attribute of graphene in the present context. In fact, room-temperature ballistic transport over distances exceeding 1 \( \mu \text{m} \) has been reported with graphene samples deposited on thin exfoliated films of h-BN \[11\], which can realistically be incorporated in the geometry of figure 1. It should also be noted here that the corrugation periods under consideration are much larger than the graphene interatomic separation; as a result, the corrugation is not expected to cause any additional mobility degradation, as opposed to the naturally occurring wrinkles in graphene which involve much shorter length scales \[7\].

The circles in figure 4(a) represent the calculated frequency of peak emission (i.e., the main radiation frequency of the \( \theta = 0^\circ \) sinusoidal trajectory) plotted as a function of corrugation period \( \Lambda \) for fixed amplitude \( A = 50 \text{ nm} \). The observed red-shift with increasing period is in agreement with equation (1), although the values plotted in figure 4(a) are consistently smaller than the predictions of this simple formula. The discrepancy is again due to the underlying assumption of equation (1) (i.e., \( A \ll \Lambda \)), which is not fully justified in the geometries under study. As shown in figure 4(b) (where a fixed period of 300 nm is assumed), the emission frequency also decreases with increasing corrugation amplitude, which can be ascribed to the resulting increase in travel time per period. In any case, the important conclusion from these plots is that radiation frequencies across a wide portion of the THz spectrum can be obtained with realistic values of \( A \) and \( \Lambda \) on the order of several ten and a few hundred nanometers, respectively. Such values can be readily obtained in surface-relief gratings based on current nanofabrication technologies, which can then be used as substrates for the conformal adhesion of pre-synthesized graphene sheets.

The same FDTD simulations can also be used to determine the far-field radiation pattern of the emitted light. Exemplary results are shown in figure 5, for the specific
Figure 4. Peak emission frequency (circles) and total output power per unit area (squares) of corrugated graphene versus corrugation period (a) and versus corrugation amplitude (b). The amplitude and period in (a) and (b) are fixed at 50 nm and 300 nm, respectively. The output power was computed using equation (6) assuming a carrier density of $5 \times 10^{12}$ cm$^{-2}$ and room temperature.

Figure 5. Far-field radiation pattern of an electron traveling along the $\theta = 0^\circ$ trajectory on a corrugated sheet of graphene with 300-nm period and 50-nm amplitude. Panels (a) and (b) contain polar plots of the emitted light intensity at the frequency of peak emission on the $xy$ and $xz$ planes, respectively (defined relative to the system of coordinates of figure 1).

In the case of an electron traveling along the $\theta = 0^\circ$ trajectory on a sinusoidal corrugation with 300-nm period and 50-nm amplitude, each panel in the figure contains a polar plot of the far-field output intensity (at the frequency of peak emission), on a different Cartesian plane relative to the system of coordinates of figure 1. In the limit of negligibly small corrugation amplitude relative to the period, simple electromagnetic considerations based on the Larmor formula indicate that the far-field radiation pattern approaches that of a dipole oriented along the $z$ direction. The results plotted in figure 5 still resemble such a field profile, except that they are not symmetric with respect to reflections about the $yz$ plane, with more light emitted along the negative $x$ direction (i.e., antiparallel to the electronic motion) compared to the positive $x$ direction. This shape is related to the $x$-dependence of the phase of the equivalent polarization density of equation (5). From a practical standpoint, the data of figure 5 suggest that the experimental measurement of the radiated light will require either a side-emission configuration or additional optical elements for vertical outcoupling.

Finally, we address the question of how much THz power can be obtained based on the radiation mechanism described in this work. Specifically, we consider an n-doped corrugated graphene sheet of electron density $N$, under the action of an externally applied voltage producing an electric field $\mathbf{E}$ in the direction of the corrugation. The total spontaneous emission power spectrum $P_{\text{tot}}(\nu)$ can be computed as follows:

$$P_{\text{tot}}(\nu) = 4 \sum_{\mathbf{k}} P_{e}(\mathbf{k}, \nu) f(\mathbf{k}) \left[ 1 - f(\mathbf{k}') \right],$$

where the factor of 4 accounts for the spin and valley degeneracies, $P_{e}(\mathbf{k}, \nu)$ is the power radiated by an electron in the conduction-band Bloch state of wavevector $\mathbf{k}$, $f(\mathbf{k})$ is the probability that this initial state is occupied, and $1 - f(\mathbf{k}')$ is the probability that the corresponding final state after photon emission is empty. The wavevector $\mathbf{k}'$ of the latter state is related to the initial electronic wavevector $\mathbf{k}$ by the requirements of conservation of energy, i.e., $E(\mathbf{k}) - E(\mathbf{k}') = \hbar \nu$, and conservation of momentum along the $y$ direction (where the system under study has translational invariance), i.e., $k'_y = k_y$.

The single-electron radiation spectra $P_{e}(\mathbf{k}, \nu)$ for different wavevectors $\mathbf{k}$ can be calculated via the FDTD simulations described earlier (e.g., as plotted in figure 3(a)), followed by a convolution with a Lorentzian lineshape function to account for collision broadening as already discussed. For any given corrugation geometry, these spectra only depend on the electron velocity $v_{\mathbf{k}} = v_{\mathbf{F}} \mathbf{k}$, which in turn depends on the direction of $\mathbf{k}$ but not on its magnitude. Therefore, we can write $P_{e}(\mathbf{k}, \nu)$ as $P_{e}(\theta, \nu)$, where $\theta$ is
the angle between $k$ and the projection of the $x$-axis on the corrugated surface, as defined previously. To determine the occupation probabilities $f(k)$, we use a semiclassical model of carrier dynamics in crystalline solids [31], where an externally applied electric field $\mathbf{E}$ has the effect of simply displacing the equilibrium Fermi–Dirac distribution function $f_0(k)$ as a whole by the amount

$$\delta k = -\frac{q}{\hbar} \tau_F \mathbf{E} = -\frac{\mu \mathbf{F}}{v_F},$$

leading to $f(k) = f_0(k - \delta k)$. In equation (7), $\tau_k$ is the momentum relaxation time, $\mu$ is the mobility, and the second equality is obtained by assuming $\tau_k = (\mu \hbar)/(qv_F)$, an expression that appears to be consistent with the key electronic transport properties of graphene [7].

The emission spectra plotted in figure 3(b) were calculated using the procedure just described. Equation (6) can also be used to compute the total radiated power, by integrating $P_{\text{tot}}(\nu)$ over all frequencies. The squares in figures 4(a) and (b) correspond to the total power per unit sample area associated with the fundamental harmonic of the output radiation (e.g., the lower-frequency peak of figure 3(b)), plotted as a function of corrugation period $\Lambda$ and amplitude $A$, respectively. In all these calculations we used a reasonably large electron density $N$ of $5 \times 10^{12}$ cm$^{-2}$, which is well within the range of experimentally accessible values based on capacitive doping. The temperature used to compute the distribution functions was taken to be 300 K. Finally, for the effective drift velocity $\mu F$ we used a value of $2 \times 10^7$ cm s$^{-1}$, based on a recent measurement of saturation velocity in graphene at a carrier density of a few $10^{12}$ cm$^{-2}$ at room temperature [32].

If we assume a relatively small sample area of $100 \times 100$ $\mu$m$^2$, which can be readily produced using epitaxial techniques or chemical vapor deposition, the plots of figure 4 indicate that power levels of several ten nanowatt can be obtained, which are already appreciable for this frequency range. In fact, high-quality graphene samples with linear dimensions of several millimeters can now be synthesized, e.g., see [20], in which case very large spontaneous emission powers approaching a milliwatt become feasible. In this respect it is important to note that the radiation mechanism described in this work is not affected by capacitive or transit-time effects, which would otherwise limit the usable device area as in the case of other THz sources such as photomixers. Figure 4(a) also shows that the output power of corrugated graphene increases monotonically with decreasing corrugation period, which is due to the resulting increase in the curvature of the carrier trajectories. Furthermore, as illustrated in figure 4(b), for fixed period there is an optimal amplitude that maximizes the (fundamental harmonic) output power. At larger amplitudes, the higher-order harmonics (not included in these plots) are amplified instead at the expense of the lower-frequency peak. We also note that the power levels plotted in figure 4 are of the same order of magnitude, but consistently smaller, than predictions based on the approximate formula of equation (2).

From a technological standpoint, an important property that emerges from the two traces of figure 4(a) is the increase in output power with increasing emission frequency. As a result, the radiation mechanism described in this work may allow extending the frequency range of existing room-temperature THz sources such as microwave frequency multipliers and photomixers, whose output power near 1 THz can reach a few microwatt but then rapidly scales down with increasing frequency [3]. Furthermore, cyclotron-like radiation in corrugated graphene can also in principle be used to provide stimulated emission and optical gain for the development of graphene-based THz lasers. Specifically, radiation propagating in the direction of the corrugation can produce a bunching of the electron beam (via a resonant electromagnetic interaction), leading to amplification and coherent light emission, as in the case of traditional free-electron lasers [26]. Similar laser sources can therefore be envisioned based on corrugated graphene structures, combined with an optical cavity where the emitted light is reflected back and forth in the direction of the applied voltage. For a comparison with existing THz quantum cascade lasers, we note that at cryogenic temperatures the total spontaneous emission power of their active materials (in the absence of optical feedback leading to stimulated emission) is also a few ten nW for sample areas of order $100 \times 100$ $\mu$m$^2$ [33–35]. However, as the temperature is increased, their radiative efficiency rapidly decreases due to phonon-assisted nonradiative transitions between the active subbands [1], so that no detectable radiation is produced at room temperature. In the corrugated graphene emitters considered in this work, increasing the temperature mainly leads to a decrease in the maximum achievable drift velocity $\mu F$ (an effect that can be quite modest as shown in [32]), and to enhanced collision broadening of the radiation spectra. As shown in figure 3(b), the latter effect can also be relatively weak in high-quality graphene samples (particularly when deposited on h-BN), which feature room-temperature mean free paths $d$ of several hundred nanometers. As a result, if graphene-based THz ‘free-electron’ lasers can be developed, they may potentially allow for room-temperature operation.

Finally, from a fundamental standpoint, we note that the same favorable properties of graphene may also be exploited to demonstrate other mechanisms of electron-beam radiation, such as the Smith–Purcell [23, 24] and Cherenkov [25] effects, which would involve the use of planar graphene sheets in the vicinity of nanoscale gratings or photonic crystals. Altogether, these radiation phenomena represent a new paradigm for light emission in compact solid-state systems, and may find important applications in the areas of THz electronics and nanophotonics.

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