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Transformation of surface plasmon polaritons to radiation in graphene in terahertz regime

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Abstract

We demonstrate a concept that allows direct excitation of surface plasmon polaritons (SPPs) by a moving electron bunch above a single layer graphene sheet deposited on a dielectric substrate without any additional coupling requirements. We show that if the two-dimensional current in the graphene is dominated by the third order nonlinear effect when the surface electric field exceeds a moderate strength of $\sim 5\text{kV/cm}$, the SPP mode can cross the light line although the group velocity remains much smaller than the speed of light. This effect gives rise to direct transformation of SPPs into radiation. The underlying mechanism of the crossing of the SPP dispersion into the light line is the energy shift of charged particles in the nonlinear regime and the finite transport scattering time in graphene. Both the energy and lifetime of the SPPs increase with the field intensity. The radiation intensity and frequency can be tuned with an AC bias.

Keywords

terahertz, graphene, radiation, polaritons, regime, plasmon, transformation, surface

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Transformation of surface plasmon polaritons to radiation in graphene in terahertz regime

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We demonstrate a concept that allows direct excitation of surface plasmon polaritons (SPPs) by a moving electron bunch above a single layer graphene sheet deposited on a dielectric substrate without any additional coupling requirements. We show that if the two-dimensional current in the graphene is dominated by the third order nonlinear effect when the surface electric field exceeds a moderate strength of ~ 5 kV/cm, the SPP mode can cross the light line although the group velocity remains much smaller than the speed of light. This effect gives rise to direct transformation of SPPs into radiation. The underlying mechanism of the crossing of the SPP dispersion into the light line is the energy shift of charged particles in the nonlinear regime and the finite transport scattering time in graphene. Both the energy and lifetime of the SPPs increase with the field intensity. The radiation intensity and frequency can be tuned with an AC bias. © 2015 AIP Publishing LLC. [<http://dx.doi.org/10.1063/1.4922261>]

Graphene has some exceptional electronic, mechanical, and optical properties.^{1–5} Because the plasmon frequency in graphene lies in the terahertz frequency regime and is found to be tunable, intense research has been carried out to search for graphene based radiation sources, detectors, and modulators,^{6–11} especially in the terahertz region. Recent experimental and theoretical works have shown that graphene can support surface plasmon polaritons (SPPs).^{12–17} Due to the unique properties of graphene including high mobility and universal conductance, there can be a significant enhancement of light radiation from graphene SPPs. The enhancement can be up to two orders of magnitude when a graphene sheet is deposited on a periodic dielectric substrate. Graphene SPPs are particularly useful at low frequency as the plasmon frequency of graphene with favourable charge carrier concentration lies in the 1–50 THz frequency regime. This makes graphene an ideal candidate for developing low frequency light radiation sources such as terahertz emitters. For a uniform graphene structure, the SPP modes lie below the light dispersion line of the dielectrics. Therefore, the direct conversion of SPPs to light is impossible. The key challenge in graphene photonics and plasmonics is enabling SPPs to be excited by an incident field. The two most commonly used techniques involve using a prism to allow for wavevector matching, or introducing a periodic microstructure either directly in the graphene sheet or in the substrate dielectrics to cause band folding. The band folding of the SPP dispersion under a periodic potential brings the SPP mode at large q back into the first Brillouin zone. Thus, the SPP mode crosses the light line, and SPP-light coupling becomes possible.¹⁸ This technique requires precise control of the periodicity.

Recently, Bludov *et al.*¹⁹ considered nonlinear transverse electric (TE)-polarized surface polaritons on graphene. They

included the nonlinear dielectric function of the substrate in the SPP calculation and found that the system can support and stabilize nonlinear TE plasmon polaritons. The results rely on a strong nonlinear dielectric as a substrate. On the other hand, it has been shown that graphene is a strong nonlinear material.^{20–24} The required threshold field for graphene to exhibit the nonlinear effect is rather moderate in the order of a few kV/cm. Gorbach recently derived an amplitude equation for nonlinear transverse magnetic (TM) and TE surface plasmon waves supported by graphene.²⁵ Significant nonlinear optical interactions at the few photon levels in graphene have been realized.²⁶ Therefore, it is reasonable to believe that with a certain substrate, graphene will enter the nonlinear regime first, while the dielectric substrate remains as a linear system.

In this letter, we demonstrate a mechanism that can bring the SPP mode above the light line without the requirement of structural modification. It has been shown that under a moderate electric field strength of a few kV/cm, the third order conductivity in graphene exceeds the linear conductivity. Both the interband and intraband conductivities depend strongly on the amplitude and frequency of the electrical field. We shall show that the combination of the nonlinear surface current and a finite relaxation time can lift the long wavelength SPP mode above the light line of the dielectric, making the SPP-radiation transformation immediately possible in the absence of any other zone-folding or diffraction mechanism. This result is of significance in understanding the high field electronic transport and SPP excitation in graphene.

Let us consider a graphene sheet grown on dielectrics with a dielectric constant ϵ_2 . Above the graphene is a different type of dielectric with dielectric constant ϵ_1 . The model system is shown schematically in Fig. 1. The intraband conductivity can be obtained by solving the Boltzmann equation

FIG. 1. The structure considered when calculation SPP dispersion relations. A layer of graphene is sandwiched between two dielectrics.

$$\frac{\partial f}{\partial t} - e\mathbf{E} \cdot \frac{\partial f}{\partial \mathbf{p}} = -\frac{f - f_0}{\tau}. \quad (1)$$

Here, f is the distribution function in the presence of an applied field, f_0 is the equilibrium distribution function, which is taken to be the Fermi-Dirac distribution function, $f_0^\pm = \left[1 + \exp\left(\frac{\pm v_F p - \mu}{k_B T}\right)\right]^{-1}$. Here, $\pm$ indicates the conduction and valence band, μ is the chemical potential, and T is the absolute temperature. The linear conductivity is given as

$$\begin{aligned} \sigma_1 &= \frac{4\sigma_0}{\pi} \frac{1}{\gamma - i\omega} \left\{ \mu + 2k_B T \ln \left[1 + \exp\left(-\frac{\mu}{k_B T}\right) \right] \right\} \\ &= \frac{4\sigma_0}{\pi} \frac{1}{\gamma - i\omega} I_1, \end{aligned} \quad (2)$$

where $\sigma_0 = e^2/\hbar$ and $\gamma = \hbar/\tau$ is the relaxation rate. The second order current is zero in uniform systems. The third order conductivity can be written as

$$\sigma_3 = \sigma_1 \frac{3}{4} \phi^2 E_0^2 \frac{(\gamma^2 - 6\omega^2) + 5i\gamma\omega}{(\gamma^2 - 6\omega^2)^2 + 25\gamma^2\omega^2} \frac{I_3}{I_1}, \quad (3)$$

where $\phi = e\hbar v_F \cong 2 \times 10^8$ eV² (cm/V), E_0 is the amplitude of the electric field on graphene, and

$$I_3 = \int_0^\infty dp \left(\frac{1}{p} \frac{\partial f_0}{\partial p} - \frac{\partial^2 f_0}{\partial p^2} - p \frac{\partial^3 f_0}{\partial p^3} \right),$$

where $f_0 = f_0^+ - f_0^-$. Fig. 2 shows the frequency dependent conductivity ratio σ_3/σ_1 in the 0–2 THz region. The nonlinear effect increases with decreasing chemical potential and initially increases then decreases as a function of frequency (cf. Fig. 2). A negative real component of the conductivity has been observed in other graphene structures in both the linear and nonlinear regimes. This property has been proposed for amplification of terahertz radiation.^{22,27,28} The most interesting nonlinear effect is that the imaginary part for σ_3 is entirely negative. This makes the excitation of a TE SPP mode in graphene possible. We found that the negative imaginary part of the conductivity is necessary for the TM SPP mode to cross the light line. For a field amplitude greater than 5 kV/cm, the conductivity is dominated by the nonlinear term.

By using Maxwell's equations and the appropriate boundary conditions, we arrive at the general dispersion relationships for the TM and TE modes

FIG. 2. The ratio of nonlinear to linear intraband conductivity as a function of frequency.

$$\frac{\epsilon_1}{\kappa_1} + \frac{\epsilon_2}{\kappa_2} + \frac{4i\alpha}{\omega} \tilde{\sigma} = 0, \quad (4)$$

$$\kappa_1 + \kappa_2 - 4i\alpha\omega\tilde{\sigma} = 0, \quad (5)$$

where Eqs. (4) and (5) are the TM and TE modes, respectively, and $\tilde{\sigma}$ is the total conductivity of graphene. Here, ω is the frequency of SPPs in units of energy. For low energy SPPs in doped graphene, only the intraband transitions contribute to the optical process.

In this work, we only consider the TM SPP mode. Assuming the dielectric above graphene is air, $\epsilon_1 = 1$. In the linear regime, in the absence of relaxation, Eq. (4) is a real equation and the solution of $\omega(q)$ always lies outside the light line. In the more realistic case of finite relaxation time, Eq. (4) becomes complex. Depending on the excitation method, one can either choose a complex ω and a real q , or a real ω and a complex q . In the case of complex ω , a finite relaxation rate pushes the SPP mode further away from the light line. In the case of complex q , the SPP mode can apparently cross the light line. While this seems to suggest that SPPs could potentially be excited in this region, realistically this is not possible as the imaginary part of q is considerably larger than the real component, meaning that the mode is overdamped. Propagation therefore would not occur.

To make our analysis simpler, we only consider the linear regime and strong nonlinear regime, where the conductivity is completely dominated by the nonlinear term. In the linear regime, the SPP mode is determined by Eq. (4), where $\sigma = \sigma_1$. The SPP mode in the presence of finite relaxation is shown in Fig. 3. The effect of the relaxation rate is to shift the SPP mode further away from the light line. In the strong nonlinear regime, σ_3 dominates the current, which oscillates at triple frequency 3ω . We therefore need to match the field components with frequency 3ω either side of the boundary and use σ_3 as the surface conductivity term. This gives the nonlinear SPP mode as

$$\frac{\epsilon_1}{\kappa_1(3\omega)} + \frac{\epsilon_2}{\kappa_2(3\omega)} + \frac{4i\alpha}{3\omega} \sigma_3(3\omega) = 0. \quad (6)$$

FIG. 3. Dispersion relation of SPPs in the linear limit (dotted-dashed line) and in the strong nonlinear limit (solid pink line). The two black lines are the light lines in dielectrics 1 and 2. The electron beam velocity is $0.42c$, as shown by the red straight line. Excitation of SPPs by an electron beam is determined by the intersection of the beam line and the SPP line, or the working point. The temperature is 300 K , the chemical potential $\mu = 50\text{ meV}$, and the relaxation rate is $\gamma = 2.6\text{ meV}$.

In Fig. 3, the nonlinear SPP mode for complex ω is shown. Compared to the linear mode, the dispersion of the nonlinear mode shifts vertically. Due to this vertical shift, the dispersion relations cross into the light cone. For the case of complex q , we found that increasing the electric field causes the dispersion relation to also lie within the light cone; however, increasing the field also increases the damping of the SPPs, making this method of excitation completely infeasible. For complex ω , however, not only does the finite electric field allow the dispersion to lie within the light cone but also the decay rate decreases with increasing electric field. For sufficiently high fields ($E \geq 5\text{ kV/cm}$), the ratio $\text{Im}(\omega)/\text{Re}(\omega) < 1$ as can be seen in Fig. 4. In the linear regime, apart from the fact that the SPP mode is far away from the light line, this ratio is much greater than 1 in the small q region. In the

FIG. 4. Ratio of imaginary part of the SPP mode to the real part of the SPP mode. In the strongly nonlinear region, the ratio is less than 1 and propagation of SPPs is possible.

nonlinear regime, the ratio can be smaller than 1 in the small q region, where the SPP mode is above the light line.

These results paint an interesting picture about the effect of the electric field on the dispersion relations. Despite the inclusion of a finite relaxation rate causing the dispersion curve to deviate from the light line in the low frequency limit, an increase in the incident field is found to fix this and allow for excitation. When the relaxation rate is not included, increasing the field does not cause the SPP dispersion curve to cross into the light cones, the effect only occurs when a finite relaxation rate is included. Our finding here is that the interplay of the nonlinear conductivity and the finite relaxation can give rise to the unexpected possibility that SPPs can be directly transformed into radiation inside the light cone without any additional structural modification for momentum matching.

We shall now consider excitation of the nonlinear SPPs by an electron bunch moving with velocity u_0 above the graphene sheet. This corresponds to an excitation at the working point where the electron beam line intersects with the nonlinear SPP dispersion relation (cf. Fig. 3). The components of the electromagnetic field of the bunch oscillating at 3ω are given as

$$E_z = -\frac{e\kappa_c}{6\omega\epsilon_0} e^{i\kappa_c(y-y_0)} e^{ik_z z}, \quad (7)$$

$$H_x = -\frac{q}{2} e^{i\kappa_c(y-y_0)} e^{ik_z z}, \quad (8)$$

where $\kappa_c = \sqrt{k_0^2 - k_z^2}$, $k_z = 3\omega/u_0$ and y_0 is the position of the electron bunch. By using the boundary conditions at $y=0$,

$$E_z^I + E_z^R = E_z^H, \quad H_x^H - (H_x^I + H_x^R) = \sigma_3(E_z^H + E_0). \quad (9)$$

In the boundary equation, the H field discontinuity is determined by the total current $j(3\omega)$, which is in response to the total surface field $(E_z^H + E_0)$. The total field is the linear combination a large AC bias field E_0 and the emitted field. The emitted field is much smaller than E_0 for the entire frequency region, except at resonance. Therefore, we shall only retain the E_0 dependence in the nonlinear conductivity. We obtain the emitted field amplitude

$$A_1 = \frac{\left(\frac{i\epsilon_2\kappa_c}{\kappa_2} + 1\right) \frac{e}{2} e^{i\kappa_c y_0} + \sigma_3 E_0}{\frac{i3\omega\epsilon_0\epsilon_1}{\kappa_1} + \frac{i3\omega\epsilon_0\epsilon_2}{\kappa_2} - \sigma_3}, \quad (10)$$

$$A_2 = A_1 - \frac{e\kappa_c}{6\omega\epsilon_0} e^{i\kappa_c y_0}.$$

Fig. 5 shows that the radiation field is directly transformed from the nonlinear SPP. The radiation spectrum is a sharply peaked function centered around the frequency as determined by the working point (the intersection) of the electron beam and the SPP dispersion. For the beam speed of $0.42c$, the peak frequency is 0.39 THz . The peak height is predominantly determined by the nonlinear field E_0 , as the second term in the numerator in A_1 is much stronger than the field of the probing electron beam. The radiation amplitude

FIG. 5. Radiation field distribution due to the nonlinear SPP in dielectrics, the peak frequency is 0.39 THz, $T = 300$ K, and $\mu = 50$ meV. The field amplitude driving the nonlinear current is 6500 V/cm.

in the vicinity of the graphene sheet exceeds the probing field by many orders of magnitude. The peak frequency, which is determined by the root of the denominator of A_1 , is also strongly dependent on E_0 via the nonlinear conductivity. Increasing E_0 , results in a higher SPP dispersion and a higher working point. This result suggests that the nonlinear current in graphene can be used as the tuning parameter for the radiation field. In addition, the frequency can also be tuned by the electron beam. Unlike the case of graphene on a periodic structure where radiation frequency increases with beam speed, here the radiation frequency decreases with the beam speed. A high beam speed results in a lower working point, since the SPP energy decreases with q . In the case of graphene on a periodic structure, the SPP energy in the second band after zone folding increases as q decreases. The peak height was also found to increase with the beam energy.

It should be pointed out that although direct transformation of SPPs to radiation is possible in the current model, the propagation of the SPP generated radiation is restricted by the heavy attenuation along the direction perpendicular to the interface. The propagation in the dielectric side is more efficient than that on the air side. The radiation is strongly localized near the interface. Therefore, it may not be practical to use this mechanism for developing emitters. To reduce the attenuation, a stronger field would be required.

In conclusion, we have shown that it is possible for the SPP dispersion to cross the light line in a graphene on top of dielectrics. Due to a nonlinear current in the graphene, the SPP mode has a higher phase velocity than the phase velocity of the light in the dielectrics. When excited by an electron beam moving at speed faster than the light in the dielectrics, the SPP can be transformed into radiation.

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- ¹A. H. C. Neto, F. Guinea, N. M. R. Peres, K. S. Novoselov, and A. K. Geim, *Rev. Mod. Phys.* **81**, 109 (2009).
- ²K. S. Novoselov, A. K. Geim, S. Morozov, D. Jiang, M. Katsnelson, I. Grigorieva, S. Dubonos, and A. Firsov, *Nature* **438**, 197 (2005).
- ³J. Horng, C.-F. Chen, B. Geng, C. Girit, Y. Zhang, Z. Hao, H. A. Bechtel, M. Martin, A. Zettl, M. F. Crommie, Y. R. Shen, and F. Wang, *Phys. Rev. B* **83**, 165113 (2011).
- ⁴Z. Fei, A. S. Rodin, W. Gannett, S. Dai, W. Regan, M. Wagner, M. K. Liu, A. S. McLeod, G. Dominguez, M. Thiemens *et al.*, *Nat. Nanotechnol.* **8**, 821–825 (2013).
- ⁵Z. Q. Li, E. A. Henriksen, Z. Jiang, Z. Hao, M. C. Martin, P. Kim, H. L. Stormer, and D. N. Basov, *Nat. Phys.* **4**, 532–535 (2008).
- ⁶M. Liu, X. Yin, E. Ulin-Avila, B. Geng, T. Zentgraf, L. Ju, F. Wang, and X. Zhang, *Nature* **474**, 64 (2011).
- ⁷A. K. Geim, *Science* **324**, 1530 (2009).
- ⁸L. Ju, B. Geng, J. Horng, C. Girit, M. Martin, Z. Hao, H. A. Bechtel, X. Liang, A. Zettl, Y. R. Shen, and F. Wang, *Nat. Nanotechnol.* **6**, 630 (2011).
- ⁹M. Jablan, H. Buljan, and M. Soljacic, *Phys. Rev. B* **80**, 245435 (2009).
- ¹⁰A. Vakil and N. Engheta, *Science* **332**, 1291 (2011).
- ¹¹F. Bonaccorso, Z. Sun, T. Hasan, and A. C. Ferrari, *Nat. Photonics* **4**, 611 (2010).
- ¹²F. Rana, *IEEE Trans. Nanotechnol.* **7**, 91 (2008).
- ¹³B. Wang, X. Zhang, X. Yuan, and J. Teng, *Appl. Phys. Lett.* **100**, 131111 (2012).
- ¹⁴B. Wang, X. Zhang, F. J. Garcia-Vidal, X. Yuan, and J. Teng, *Phys. Rev. Lett.* **109**, 073901 (2012).
- ¹⁵M. Farhat, S. Guenneau, and H. Bagci, *Phys. Rev. Lett.* **111**, 237404 (2013).
- ¹⁶S. Liu, C. Zhang, M. Hu, X. Chen, P. Zhang, S. Gong, T. Zhao, and R. Zhong, *Appl. Phys. Lett.* **104**, 201104 (2014).
- ¹⁷T. Zhan, D. Han, X. Hu, X. Liu, S.-T. Chui, and J. Zi, *Phys. Rev. B* **89**, 245434 (2014).
- ¹⁸Y. V. Bludov, N. M. R. Peres, and M. I. Vasilevskiy, *Phys. Rev. B* **85**, 245409 (2012).
- ¹⁹Y. V. Bludov, D. A. Smirnova, Y. S. Kivshar, N. M. R. Peres, and M. I. Vasilevskiy, *Phys. Rev. B* **89**, 035406 (2014).
- ²⁰S. A. Mikhailov and K. Ziegler, *J. Phys.: Condens. Matter* **20**, 384204 (2008).
- ²¹A. R. Wright, X. G. Xu, J. C. Cao, and C. Zhang, *Appl. Phys. Lett.* **95**, 072101 (2009).
- ²²X. G. Xu, S. Sultan, C. Zhang, and J. C. Cao, *Appl. Phys. Lett.* **97**, 011907 (2010).
- ²³J. Wang, Y. Hernandez, M. Lotya, J. N. Coleman, and W. J. Blau, *Adv. Mater.* **21**, 2430 (2009).
- ²⁴E. Hendry, P. J. Hale, J. Moger, A. K. Savchenko, and S. A. Mikhailov, *Phys. Rev. Lett.* **105**, 097401 (2010).
- ²⁵A. V. Gorbach, *Phys. Rev. A* **87**, 013830 (2013).
- ²⁶M. Gullans, D. E. Chang, F. H. L. Koppens, F. J. Garcia de Abajo, and M. D. Lukin, *Phys. Rev. Lett.* **111**, 247401 (2013).
- ²⁷M. Ryzhii and V. Ryzhii, *Jpn. J. Appl. Phys., Part 2* **46**, L151 (2007).
- ²⁸V. Ryzhii and M. Ryzhii, *Phys. Rev. B* **79**, 245311 (2009).