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Graphene coated subwavelength wires: a theoretical investigation of emission and radiation properties

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ABSTRACT

This work analyzes the emission and radiation properties of a single optical emitter embedded in a graphene–coated subwavelength wire. We discuss the modifications of the spontaneous emission rate and the radiation efficiency as a function of the position and orientation of the dipole inside the wire. Our results show that these quantities can be enhanced by several orders of magnitude when the emission frequency coincides with one of the resonance frequencies of the graphene–coated wire. In particular, high–order plasmon resonances are excited when the emitter is moved from the wire center. Modifications resulting from varying the orientation of the dipole in the near field distribution and in the far field intensities are shown.

1. Introduction

Surface plasmon polariton results from the coherent coupling of photons to surface charge density oscillations [1,2]. In natural plasmonic materials such as metals, constitutive parameters such as conductivity and charge density are fixed, whereas in graphene, they can be tuned electrically or by chemical doping [3], which has a dramatic effect on its optical properties. As the linear band structure of graphene causes the plasmon mass to depend on the Fermi–level position, electrically tunable surface plasmons can be supported by graphene from microwaves to the mid–infrared regimes [4]. This has aroused new interest in studying graphene in the context of optical and plasmonic applications, including graphene quantum dots as a new generation of light–emitting devices [5], sensing [6], solar cells [7], to mention just a few.

Surface plasmons polaritons can be roughly divided into two categories: surface plasmon (SPs) propagating along waveguiding structures, such as an infinite flat graphene monolayer, and localized surface plasmons (LSPs) supported by spatially limited structures, such as scattering particles. Both kind of plasmon modes can be excited when a single optical emitter (such as quantum dots and single molecules) is placed next to one of the aforementioned structures. As a result of the high light confinement, an enhanced decay rate of the emitter into the plasmonic mode via the Purcell effect takes place [8]. In fact, it has recently been shown that the interplay between an optical emitter and a graphene–coated sphere leads to an enhancement of the spontaneous emission into confined plasmonic modes [9]. A variety of structures such as infinite graphene monolayers [10–13], ribbons or nanometer sized disks [14] and double graphene waveguides [15,16] have been the object of intensive research over the last few years exploring the possibility to engineer surface plasmon mode density of states to control emission properties.

In this paper we consider a cylindrical dielectric core coated with a graphene layer and we investigate the role of LSPs in modifying the emission and radiation rates of an emitter placed inside the graphene–coated cylinder. In this context, several works have focused on the influence that the eigenmodes play on circular dielectric and metallic waveguides [17–19] or on single wall carbon nanotubes [20]. Such systems offer a new tool for the interplay between light and matter at nanometer scale [21,22]. Other works have presented two dimensional calculations extending such studies to the role played by the shape of the waveguide [23,24]. A two–dimensional setting has important properties that allow a better understanding of the three dimensional case [17,28], since qualitatively similar trends are held in three dimensions [24]. Experimental advances in the manufacturing of micro–structures based on insulating materials such as SiO_2 doped with molecular ions [25], the possibility of encapsulating single atoms, molecules and compounds into graphene wires [26] and the fact that, as a result of the van der Waals force, a graphene sheet can be tightly

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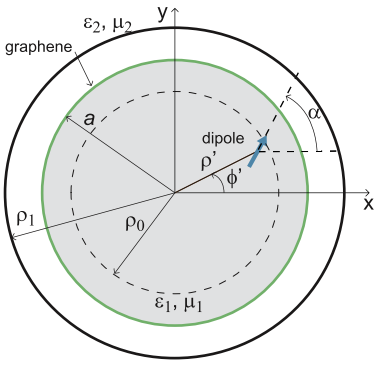

Fig. 1. Schematic illustration of the system. An optical dipole emitter is inside a graphene–coated dielectric cylinder. The cylinder radius is a and the graphene surface conductivity is σ.

coated on a fiber surface [27], encourages further research into emission and radiation characteristics of such graphene based systems.

This paper is organized as follows. First, in Section 2 we sketch an analytical method based on the separation of variables approach and obtain a solution for the electromagnetic field scattered by a graphene–coated wire when an oscillating line dipole source is located at an arbitrary position inside the wire cylinder. We derive analytical expressions for the power emitted and radiated by the source. In Section 3 we present examples of emission and radiation decay rates corresponding to wires tightly coated with a graphene layer and compare the results with those obtained in the absence of the graphene coating. Finally, concluding remarks are provided in Section 4. The Gaussian system of units is used and an $\exp(-i\omega t)$ time–dependence is implicit throughout the paper, with ω as the angular frequency, t as the time, and $i=\sqrt{-1}$. The symbols Re and Im are respectively used for denoting the real and imaginary parts of a complex quantity.

2. Theory

We consider a graphene coated cylinder with circular cross section (radius a) centered at $x=0$, $y=0$ (Fig. 1). The wire substrate is characterized by the electric permittivity ε_1 and the magnetic permeability μ_1. The coated wire is embedded in a transparent medium with electric permittivity ε_2 and magnetic permeability μ_2. We assume that radius a is sufficiently large as to describe the optical properties of the wire as those of a circular cylinder characterized by the same surface conductivity as planar graphene [20]. In this way, the graphene layer is considered as an infinitesimally thin, local and isotropic two–sided layer with frequency–dependent surface conductivity $\sigma(\omega)$ given by the Kubo formula [29,30], which can be read as $\sigma=\sigma^{intra}+\sigma^{inter}$, with the intraband and interband contributions being

$$\sigma^{intra}(\omega)=\frac{2ie^2k_BT}{\pi\hbar(\omega+i\gamma_c)}\ln[2\cosh(\mu_c/2k_BT)], \tag{1}$$

$$\sigma^{inter}(\omega)=\frac{e^2}{\hbar}\left\{\frac{1}{2}+\frac{1}{\pi}\arctan[(\omega-2\mu_c)/2k_BT]-\frac{i}{2\pi}\ln\left[\frac{(\omega+2\mu_c)^2}{(\omega-2\mu_c)^2+(2k_BT)^2}\right]\right\}, \tag{2}$$

where μ_c is the chemical potential (controlled with the help of a gate voltage), γ_c the carriers scattering rate, e the electron charge, k_B the Boltzmann constant and $\hbar$ the reduced Planck constant. The intraband contribution dominates for large doping $\mu_c >> k_BT$ and it is a generalization of the Drude model for the case of arbitrary band structure, whereas the interband contribution dominates for large frequencies $\hbar\omega \geq \mu_c$. A line dipole source (whose axis lies along the $\hat{z}$ axis) with a dipole moment $\vec{p}=p(\cos\alpha\,\hat{x}+\sin\alpha\,\hat{y})$ is placed inside the cylinder, at position $\vec{r}'=\rho'\hat{r}+\phi'\hat{\phi}$ ($\rho'<a$). The dipole is aligned at an angle α with respect to the $\hat{x}$ axis, as indicated in Fig. 1. The current density of the electric dipole is

$$\vec{j}(\vec{r})=-i\omega\vec{p}\,\delta(\vec{r}-\vec{r}')=-i\omega\vec{p}\,\frac{1}{\rho}\delta(\rho-\rho')\delta(\phi-\phi'). \tag{3}$$

In an unbounded medium, the dipole fields are obtained from the vector potential $\vec{A}$ (refer to Appendix A for its derivation),

$$\vec{A}(\rho,\phi)=\sum_{m=-\infty}^{+\infty}\pi k_0 J_m(k_1\rho_<)H_m^{(1)}(k_1\rho_>)\,e^{im(\phi-\phi')}\times\left[p_\rho\hat{r}+p_\phi\hat{\phi}\right], \tag{4}$$

$$\vec{H}(\rho,\phi)=\nabla\times\vec{A}=\hat{z}\varphi(\rho,\phi), \tag{5}$$

$$\vec{E}(\rho,\phi)=\frac{i}{k_0\varepsilon_1}\nabla\times\vec{H}(\rho,\phi)=-\frac{i}{k_0\varepsilon_1}\hat{z}\times\nabla_t\varphi, \tag{6}$$

where $\varphi(\rho,\phi)$ is the non–zero component of the total magnetic field along the axis of the wire ($\hat{z}$ axis), $\nabla_t=\hat{r}\frac{\partial}{\partial\rho}+\hat{\phi}\frac{1}{\rho}\frac{\partial}{\partial\phi}$ is the transverse part of the ∇ operator, $k_0=\omega/c$ is the modulus of the photon wave vector in vacuum, ω is the angular frequency, c is the vacuum speed of light, $\rho_<$ ($\rho_>$) is the smaller (larger) of ρ and ρ', p_ρ and p_ϕ are the projection of vector $\vec{p}$ on the $\hat{r}$ and $\hat{\phi}$ axis, respectively, and J_m and $H_m^{(1)}$ are the nth Bessel and Hankel functions of the first kind, respectively. From Eqs. (4) and (5), we obtain the primary magnetic field emitted by the dipole,

$$\varphi_i(\rho,\phi)=\sum_{m=-\infty}^{+\infty}\pi k_0k_1J_m(k_1\rho')\times\left[H_m^{(1)'}(k_1\rho)\,p_\theta-im\frac{H_m^{(1)}(k_1\rho)}{k_1\rho}\,p_\rho\right]e^{im(\phi-\phi')}, \tag{7}$$

for $\rho>\rho'$, and

$$\varphi_i(\rho,\phi)=\sum_{m=-\infty}^{+\infty}\pi k_0k_1H_m^{(1)}(k_1\rho')\times\left[J_m'(k_1\rho)\,p_\theta-im\frac{J_m(k_1\rho)}{k_1\rho}\,p_\rho\right]e^{im(\phi-\phi')}, \tag{8}$$

for $\rho<\rho'$.

When the dipole is located inside the coated wire cylinder, the scattered magnetic field along the axis of the wire, *i.e*, the $\hat{z}$ component, denoted by $\varphi_s^{(j)}$ ($j=1,2$), is expanded as a series of cylindrical harmonics, one for the internal region ($\rho<a$, superscript 1) and another one for the external region ($\rho>a$, superscript 2),

$$\varphi_s^{(1)}(\rho,\phi)=\sum_{m=-\infty}^{+\infty}a_mJ_m(k_1\rho)e^{im\phi}, \tag{9}$$

$$\varphi_s^{(2)}(\rho,\phi)=\sum_{m=-\infty}^{+\infty}b_mH_m^{(1)}(k_2\rho)e^{im\phi}, \tag{10}$$

where a_m and b_m are unknown complex coefficients. Due to the graphene coating, the tangential components of the magnetic field are no longer continuous across the boundary as they were in the case of uncoated cylinders. Considering this, the boundary conditions for our case can be expressed as

$$\frac{1}{\varepsilon_1}\frac{\partial}{\partial\rho}(\varphi_i+\varphi_s^{(1)})|_{\rho=a}=\frac{1}{\varepsilon_2}\frac{\partial}{\partial\rho}\varphi_s^{(2)}|_{\rho=a}, \tag{11}$$