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Enhancement, suppression of the emission and the energy transfer by using a graphene subwavelength wire

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ABSTRACT

The present work focuses on theoretically research on the spontaneous emission and the energy transfer process between two single optical emitters placed close to a graphene coated wire. The localized surface plasmons (LSPs) supported by the structure provide decay channels which lead to an enhancement of the emission and radiation decay rates as well as an improvement in the energy transfer between two dipole emitters. Modifications resulting from varying the orientation of dipole moments in these quantities are shown. We find that the radiation and the energy transfer efficiencies can be largely reduced at a specific frequency depending on the emitter location. By dynamically tuning the chemical potential of graphene coating, the spectral region where the dipole–field interaction is enhanced can be chosen over a wide range.

1. Introduction

It is known that the radiation characteristics of a quantum emitter (atom, molecule, quantum dots or wires) depends on the environment [1], and the spontaneous emission can be controlled by modifying the photonic density of states to the oscillating dipole [2]. A way to achieve this effect consists of coupling the quantum emitter to resonant optical modes, like wave guiding or surface plasmon (SP) modes, which provide new channels to selectively increase the emission rate into these modes [3–7].

The range of the quantum emitters interaction can also be enhanced by coupling these emitters to a designed wave guiding or SP mode environments. A variety of structures such as uniform planar microcavities [8], cylindrical nanowires [9] nanowaveguides [10] and metal films [11] have been the object of intensive research over the last few years due to the possibility to engineer wave guiding or SP mode distributions to mediate the energy transfer between quantum emitters.

High radiative quantum efficiency, large spontaneous emission rate and efficient light energy transfer are some of the main features to achieve practical nanophotonic devices, such as nanoscale antennas ([3,4,12] and references therein), plasmon lasers [13] and photon–plasmon coupler nanostructures [14,15].

The advent of graphene, which offers electrically tunable, low loss and highly confined SPs [16–18], has widened possibilities to light control in a wide frequency range allowing the movement of metal plasmonic applications, currently working in the visible range, to the mid–infrarred and terahertz frequencies [19–22]. Special phase matching techniques, such as prism or grating coupling, extensively used in the metallic and metamaterial case [23–26], have been considered involving graphene in Refs. [27–30]. In addition, the improvement of the spontaneous emission of a quantum emitter close to a graphene sheet and the possibility to control the strength of the interaction between two emitters placed in the vicinity of the sheet [31,32], has aroused a particular interest in studying other geometries that include graphene [33–37].

In this paper we consider a cylindrical dielectric core of circular cross section coated with a graphene layer. In this context, we have focused on the influence that the SPs play in the emission characteristics and the optical interaction between two line dipole emitters in close proximity to a graphene wire.

In a previous work [38] we investigated the emission and radiation properties of a single emitter placed inside a graphene coated cylinder. Unlike the results found in this work, in the present contribution we find radiation spectra highly depending on the orientation of the emitter dipole moment. In addition, both the resonant enhancements as well as a significant reduction in the power

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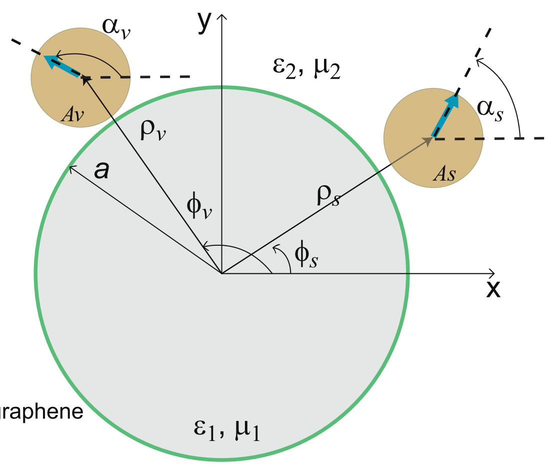

Fig. 1. Schematic illustration of the system.

radiated by the emitter have been found. Curiously, as its active media counterpart [39] the graphene coated wire has the ability to enhance the radiated power by an emitter in proximity to it as well as to reduce at near zero values this power. We also find large enhancements in the energy transfer between line dipole emitters which can be controlled by electrical tuning of the plasmon resonances.

This paper is organized as follows. First, in Section 2 we sketch the boundary–value problem for the diffraction of a wave emitted by an oscillating dipolar line source placed in close proximity to a dielectric cylinder wrapped with a graphene sheet and derive analytical expressions for the emitted and radiated power efficiencies. We then include a second dipolar line source and briefly deal with the problem of the coupled system, providing a known expression to calculate the energy transfer between two dipoles. In Section 3 we present the examples analyzing the magnitude of the radiation efficiency as well as the energy transfer decay rate. 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

2.1. Electromagnetic field scattered by a line dipolar source

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. A two dimensional emitter (a line dipole source whose axis lies along the $\hat{z}$ axis) with a dipole moment $\vec{p}_s = p(\cos\alpha\,\hat{x} + \sin\alpha\,\hat{y})$ is placed outside the cylinder, at position $\vec{r}_s = \rho_s\hat{\rho} + \phi_s\hat{\phi}$ ($\rho_s > a$). The dipole is aligned at an angle α_s with respect to the $\hat{x}$ axis, as indicated in Fig. 1. The current density of the electric dipole is

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

In an unbounded medium 2 the dipole fields are obtained from the vector potential $\vec{A}$ [38],

$$\vec{A}(\rho,\phi) = \sum_{m=-\infty}^{+\infty} \pi k_0 J_m(k_2\rho_<)H_m^{(1)}(k_2\rho_>)\,e^{im(\phi-\phi_s)} \times \left[p_\rho\hat{\rho} + p_\phi\hat{\phi}\right], \tag{2}$$

$$\vec{H}(\rho,\phi) = \nabla\times\vec{A} = \hat{z}\varphi(\rho,\phi), \tag{3}$$

$$\vec{E}(\rho,\phi) = \frac{i}{k_0\varepsilon_2}\nabla\times\vec{H}(\rho,\phi) = -\frac{i}{k_0\varepsilon_2}\hat{z}\times\nabla_t\varphi, \tag{4}$$

where $\varphi(\rho,\,\phi)$ is the non–zero component of the total magnetic field along the axis of the wire ($\hat{z}$ axis), $\nabla_t = \hat{\rho}\frac{\partial}{\partial\rho} + \hat{\phi}\frac{1}{\rho}\frac{\partial}{\partial\phi}$ 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_s, p_ρ and p_ϕ are the projection of vector $\vec{p}_s$ on the $\hat{\rho}$ 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. (2) and (3), we obtain the primary magnetic field emitted by the dipole,

$$\varphi_i(\rho,\phi) = \sum_{m=-\infty}^{+\infty} \pi k_0 k_2 J_m(k_2\rho_s) \times\left[H_m^{(1)'}(k_2\rho)\,p_\theta - im\frac{H_m^{(1)}(k_2\rho)}{k_2\rho}\,p_\rho\right]e^{im(\phi-\phi_s)}, \tag{5}$$

for $\rho > \rho_s$, and

$$\varphi_i(\rho,\phi) = \sum_{m=-\infty}^{+\infty} \pi k_0 k_2 H_m^{(1)}(k_2\rho_s) \times\left[J_m'(k_2\rho)\,p_\theta - im\frac{J_m(k_2\rho)}{k_2\rho}\,p_\rho\right]e^{im(\phi-\phi^s)}, \tag{6}$$

for $\rho < \rho_s$.

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 ($\rho < a$),

$$\varphi_{sc}^{(1)}(\rho,\phi) = \sum_{m=-\infty}^{+\infty} a_m J_m(k_1\rho)e^{im\phi}, \tag{7}$$

and another one for the external region ($\rho > a$),

$$\varphi_{sc}^{(2)}(\rho,\phi) = \sum_{m=-\infty}^{+\infty} b_m H_m^{(1)}(k_2\rho)e^{im\phi}, \tag{8}$$

where a_m and b_m are unknown complex coefficients. The boundary conditions for the graphene–coated cylinder at $\rho = a$ can be expressed as [38]

$$\frac{1}{\varepsilon_1}\frac{\partial}{\partial\rho}\varphi_{sc}^{(1)}|_{\rho=a} = \frac{1}{\varepsilon_2}\frac{\partial}{\partial\rho}(\varphi_i + \varphi_{sc}^{(2)})|_{\rho=a}, \tag{9}$$

and

$$(\varphi_i + \varphi_{sc}^{(2)})|_{\rho=a} - \varphi_{sc}^{(1)}|_{\rho=a} = \frac{4\pi\sigma}{ck_0\varepsilon_1}i\frac{\partial}{\partial\rho}\varphi_{sc}^{(1)}|_{\rho=a}. \tag{10}$$

By inserting the expressions (5), (7) and (8) into the boundary conditions, the amplitudes of the scattered field can be written as

$$a_m = f_m(y_s)\tilde{a}_m, \tag{11}$$

$$b_m = f_m(y_s)\tilde{b}_m, \tag{12}$$

where