Tunable terahertz cloaking and lasing by the optically pumped graphene wrapped on a dielectric cylinder

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Abstract

The scattering properties of a dielectric cylinder wrapped by a monolayer of graphene under an external optical pump are investigated. The THz cloaking is realized due to the strong suppression of scattering at certain range of frequencies. While at some other frequencies, the scattering can be dramatically enhanced, and a THz lasing can be achieved due to the enormous amplification of the local field by the graphene plasmon. Both the work frequencies of THz lasing and cloaking can be tuned continuously by the quasi-Fermi energy of graphene. Our proposal can pave the way to all-optical manipulation of THz lasing and cloaking based on the optically pumped graphene.

1. Introduction

Optical scattering by subwavelength structures such as nanowires and nanospheres is of great interest in nanophotonics. In the micro- and nano-scale, the cloaking and lasing are two kinds of phenomena of opposite scattering limit. The aim of cloaking is to reduce the scattering of a given object or region of space toward the zero limit, which can be realized by specific designs of metamaterials [1–4], plasmonic cloaking [5], and ‘mantle cloaking’ [6]. On the contrary, lasing in micro- and nano-scale (also called ‘spaser’ for plasmons [7]) is to enlarge the scattering toward the limit of infinity. As a counterpart of conventional laser, plasmonic nanolaser is induced by the strong interaction among the incident light, optical gain media, and plasmonic resonators [8]. A variety of subwavelength nanolasers have been designed in recent years by using plasmonic resonators and optical gain media [9–11]. The optical gain media are mostly composed of quantum dots or dye molecules, which can compensate the loss of plasmonic resonances and supply extra energy. On the other hand, the plasmonic resonance modes can provide a strength-related feedback for the gain media, leading to the anomalous amplification of light.

Graphene, being a single atomic sheet of graphite, has attracted widespread attention in the recent decade [12]. A doped graphene can support surface plasmon polaritons (SPPs) in THz/infrared frequency regime [13, 14]. However, optical loss of graphene plasmons is still an inevitable drawback for the application based on the doped graphene. Under sufficiently strong optical pumping, graphene can be a promising active plasmonic medium at THz frequencies [15]. The population inversion of Dirac fermion and the resultant negative conductivity of graphene have been experimentally demonstrated using time-resolved spectroscopy [16]. The unique properties of optically pumped graphene enable potential applications in many functional THz devices [17–24], such as meta-surfaces with tunable strong magnetic resonances [18], PT symmetric sensors [19], THz laser in dielectric Fabri–Perot resonators [20], slot waveguides [21], periodic metallic nanostructures [22], spectral singularities in PT-diffractive grating [23] and lasing terahertz metamaterials [24].

In this work, we propose a tunable THz cloaking and lasing in the system of a dielectric cylinder wrapped by an optically pumped graphene. For the TM wave incidence, the electric dipolar resonance is a dominant term. The scattering cross section can be strongly suppressed at certain range of frequencies, which can be considered
as a THz cloaking. Meanwhile, an extremely large scattering efficiency (THz lasing) can also be realized due to the strong interaction of the localized surface plasmon mode and gain amplification of graphene. Finally, the dependence of THz cloaking and lasing on the quasi-Fermi energy, relaxation time and optical parameters of cylinder are also analyzed.

2. Theory

The system under study is schematically shown in figure 1. We consider an infinitely long cylinder with a radius $R$ wrapped by a pristine graphene monolayer. The optical parameters for the cylinder and background medium are respectively $(\varepsilon_1, \mu_1)$ and $(\varepsilon_b, \mu_b)$. We consider a non-magnetic dielectric material with $\mu_1 = 1$ in the free space ($\varepsilon_b = \mu_b = 1$). Under the external optical pump, the conductivity of the pristine graphene is given as follows [15, 16, 25]:

$$\sigma(\omega) = \sigma(\omega)_{\text{intra}} + \sigma(\omega)_{\text{inter}},$$

where the subscripts represent respectively the intra- and inter-band transition process. The form of conductivity for these two kinds of transition are:

$$\sigma(\omega)_{\text{intra}} = \frac{e^2}{4\hbar} \left( \frac{2k_b T \tau}{\pi \hbar (1 - i \omega \tau)} \right) \ln \left[ 1 + \exp \left( \frac{\mu_F}{k_b T} \right) \right],$$

$$\sigma(\omega)_{\text{inter}} = \frac{e^2}{4\hbar} \left( \tanh \left( \frac{\hbar \omega - 2 \mu_F}{4k_b T} \right) \right)$$

$$- \frac{4\hbar \omega}{i \pi} \int_0^\infty \frac{G(x) - G(\hbar \omega/2)}{(\hbar \omega)^2 - 4x^2} \, dx,$$

and

$$G(x) = \frac{\sinh (x/k_b T)}{\cosh (x/k_b T) + \cosh (x/k_b T)}.$$

where $\mu_F$ is the quasi-Fermi energy, depending on the excitation of electron-hole induced by the external infrared pumping, $\tau$ is the phenomenological relaxation time of charge carriers, $e$ is the charge of electron and $\hbar$ is the reduced Plank constant, $\omega$ is the angular frequency of the incident THz wave, $k_b$ is the Boltzmann constant and $T$ is the temperature. To achieve optical gain in equations (1)–(3), the characteristic time of the emission of the optical phonon should be much shorter than the recombination time. Under this condition, the nonequilibrium electrons and holes can emit a cascade of optical phonons, and occupy the low energy states in respectively the conduction and valence bands. In this case, the heating of the system due to the nonequilibrium carriers is small, and the effective temperature $T$ is close to the lattice temperature [25, 26].

**Figure 1.** Schematic of the system under study. A dielectric cylinder is wrapped by a monolayer of graphene under infrared pump. The IR pump can be the incident waves outside the cylinder, or guided waves along the cylinder. The background is the free space. The incident THz wave is TM polarized with the magnetic field parallel to the z-axis. The right inset represents the band structure of graphene with inverted population of carriers, and the red arrow represents the interband transition induced by the infrared pump. The green arrows indicate the recombination of electrons and holes, accompanied with the emission of THz photon.
The optical properties and applications of the dielectric cylinder wrapped by the doped graphene have been investigated intensively, including the optical modulator [27], cloaking [28], biological sensing [29], super-scatterers [30], absorption enhancement [31, 32], etc. Here, we replace the doped graphene by the pristine one and an external infrared pump is introduced. Although being the form of cylindrical geometry, the conductivity of graphene can be characterized by the same surface conductivity as that of planar graphene due to the large diameter [33, 34]. As the localized surface plasmon resonance cannot be supported for the transverse electric (TE) mode, we consider transverse magnetic (TM) mode only and the magnetic field is polarized along the z (axial) axis. The incident, scattered and internal magnetic field can be expanded by the cylindrical harmonics [35]:

\[
\begin{align*}
H_{\text{inc},z} &= -H_0 \sum_{n=-\infty}^{\infty} J_n(k_0\rho) e^{in\phi}, \quad \rho > R \\
H_{\text{sc},z} &= H_0 \sum_{n=-\infty}^{\infty} a_n H_n^{(1)}(k_0\rho) e^{in\phi}, \quad \rho > R \\
H_{\text{int},z} &= H_0 \sum_{n=-\infty}^{\infty} c_n J_n(k_0\rho) e^{in\phi}, \quad \rho < R
\end{align*}
\]

where \(H_0\) is the magnetic-field amplitude of the incident wave, \(k_0 = \omega/c\), \(k_1 = \sqrt{\varepsilon_1} \omega/c\) are respectively the wavenumber in the background and dielectric medium; \(\phi\) is the azimuthal angle, \(\rho\) is the radial component; \(J_n\) and \(H_n^{(1)}\) are respectively the \(n\)-th order Bessel and Hankel function of the first kind [30], \(a_n\) is the scattering coefficient and \(c_n\) is the internal coefficients. The scattering coefficient \(a_n\) is associated with the \(n\)-th order electric multipolar moments. The corresponding electric fields can be given as well according to \(E = \frac{j \omega}{c} \nabla \times \mathbf{H} \) [35].

The optically pumped graphene is modeled as a thin conducting film with a surface current \(J_f\). The boundary conditions at \(\rho = R\) for this scattering problem can be expressed as: \((E_{\text{inc},z} + E_{\text{sc},z}) - E_{\text{int},z} = 0\) and \((H_{\text{inc},z} + H_{\text{sc},z}) - H_{\text{int},z} = 0\), where the surface current satisfies the Ohm’s law: \(J_f = \sigma(\omega) E_{\text{inc},z}\). Thus, the scattering coefficient can be derived analytically as follows:

\[
a_n = \frac{m f_n(mx) f_n^{(1)}(x) - f_n(mx) f_n^{(1)}(x) + i \sigma f_n(mx) H_n^{(1)}(x) + i \sigma f_n(mx) H_n^{(1)}(x)}{m f_n(mx) H_n^{(1)}(x) - f_n(mx) H_n^{(1)}(x) + i \sigma f_n(mx) H_n^{(1)}(x)},
\]

where \(x = k_0R\) is the size parameter, \(m = \sqrt{\varepsilon_1/\varepsilon_0}\) is the contrast of refraction indices, and \(\sigma = \sigma(\omega)/(\varepsilon_0 c)\) is dimensionless conductivity. For \(\sigma = 0\), \(a_n\) is reduced to the case of a bare dielectric cylinder [35]. According to the obtained scattering coefficients, the extinction cross section \(C_{\text{ext}}\), scattering cross section \(C_{\text{sc}}\) and absorption cross section \(C_{\text{abs}}\) are given as follows:

\[
\begin{align*}
C_{\text{ext}} &= \frac{4}{k_0} \sum_{n=-\infty}^{\infty} \text{Re}[a_n] \\
C_{\text{sc}} &= \frac{4}{k_0} \sum_{n=-\infty}^{\infty} |a_n|^2 \\
C_{\text{abs}} &= \frac{4}{k_0} \sum_{n=-\infty}^{\infty} (\text{Re}[a_n] - |a_n|^2)
\end{align*}
\]

In the Rayleigh limit \(R \ll \lambda\), the electric dipolar moment for \(n = 1\) is the dominant term and its analytical form in equation (6) is simply given as follows:

\[
a_1 = -\frac{i \sigma}{4} (k_0 R)^2 \frac{\varepsilon_1 - 1 + i \sigma/(k_0 R)}{\varepsilon_1 + 1 + i \sigma/(k_0 R)}.
\]

Therefore, the cloaking condition for the channel \(n = 1\) due to scattering cancellation can be given as follows:

\[
\begin{align*}
\text{Re}[\varepsilon_1] - 1 - \text{Im}[\sigma]/(k_0 R) &= 0, \quad (9a) \\
\text{Im}[\varepsilon_1] + \text{Re}[\sigma]/(k_0 R) &= 0. \quad (9b)
\end{align*}
\]

Obviously, equation (9a) can be satisfied for both passive and active graphene since the imaginary part of graphene conductivity is always positive. However, if the dielectric loss of cylinder is considered (\(\text{Im}[\varepsilon_1] > 0\)), the condition for scattering cancellation in equation (9b) can be realized only when the real part of graphene conductivity become negative. As a result, the cloaking effect for dielectric cylinders wrapped by the active graphene can be better than those wrapped by the passive one when \(\text{Im}[\varepsilon_1] > 0\).

On the other hand, the resonant condition for Rayleigh scattering will be:

\[
\begin{align*}
\text{Re}[\varepsilon_1] + 1 - \text{Im}[\sigma]/(k_0 R) &= 0, \quad (10a)
\end{align*}
\]
The conductivity $\sigma(\omega)$ of optically pumped graphene as a function of quasi-Fermi energy is shown in figure 2. As is well known, $\text{Re}[\sigma]$ is always a positive quantity for a doped graphene, standing for intrinsic dissipation. For the optically pumped graphene, $\text{Re}[\sigma]$ can be positive or negative as shown in figure 2(a). The solid black line marks $\text{Re}[\sigma] = 0$, representing the transition from optical loss to gain medium as the quasi-Fermi energy increases. In the gain regime with $\text{Re}[\sigma] < 0$, the graphene plasmons can be ampliﬁed [17], which is of particular interest in this study. The corresponding $\text{Im}[\sigma]$ are also demonstrated in figure 2(b). Within the frequency regime we consider (2–10 THz), $\text{Im}[\sigma]$ increases as $\mu_F$ increases for a ﬁxed frequency, (see, e.g., the black dotted line for $f = 6$ THz), whereas $\text{Im}[\sigma]$ decreases as frequency increases (see, e.g., the black dotted line for $f = 10$ THz).

The rigorous scattering coefﬁcients derived from equation (6) are plotted in ﬁgures 2(c) and (d), where we set the parameters of cylinder as $R = 1.5 \mu m$, $\epsilon_1 = 2.1$ and $\mu_F = 0.16$ eV. At the frequency range (2 ~ 10 THz) we considered, it belongs to the Rayleigh scattering regime since $R/\lambda \ll 1$. As expected, the electric dipolar term with $n = 1$ is dominant. For a conventional cylinder without any gain, there exists a general limit [36] for scattering coefﬁcients, namely, $|a_1| \leq 1$. The limit only works for passive scatterers. For a cylinder coated by optically pumped graphene, $|a_1|$ exceeds 1 at the resonance frequency 4.76 THz as shown in ﬁgure 2(c). The abnormal resonance peak for $a_1$ can be attributed to the negativity of $\text{Re}[\sigma]$, which can amplify the localized SPPs. Except the abnormal peak, there also exists a spectrum dip at 7.9 THz. Compared to the bare cylinder (black dashed line), the scattering coefﬁcient $|a_1|$ for the dielectric cylinder coated by the optically pumped graphene has been reduced about 1 order of magnitude at 7.9 THz (the corresponding scattering cross section can be reduced by about 2 orders), which manifests a cloaking effect. The cloaking effect can become better when the loss of dielectric cylinder is introduced, for example, we set the imaginary part of permittivity $\text{Im}[\epsilon_1] = 0.06$,
optically pumped graphene plays two roles, i.e., the plasmonic resonator and gain amplification. In addition to the tunable THz cloaking, THz lasing can also be realized in this system. The calculated results indicate that the THz cloaking effect persists as the parameters of dielectric cylinder vary. The work frequency of THz cloaking undergoes a red shift as we increase the radius or the permittivity of dielectric cylinder wrapped by a doped graphene monolayer, the resonance peaks are strongly dependent on the quasi-Fermi energy. Remarkably, there exists anomalous scattering peak at the frequency of 5.17 THz for $|\alpha_2|$ is reduced by 1 order to the non-loss case. In figure 2(d), the scattering coefficients $\alpha_3$ and $\alpha_4$ are also given, which are much smaller than $\alpha_2$. There is also a resonance for $|\alpha_2|$, however, the resonance peak are small and far way the frequency regime of cloaking.

To illustrate the tunable cloaking effect, we plot the normalized scattering cross section for different quasi-Fermi energy in figure 3(a). Compared to the bare cylinder (gray dashed line), the scattering cross section for the cylinder wrapped by the optically pumped graphene can be reduced about 2 orders in magnitude at the dips of spectra. Noted that the work frequencies of THz cloaking can be tuned continuously by the quasi-Fermi energy. Interestingly, the cloaking effect is weakly dependent on the relaxation time $\tau$ as shown in figure 3(b). Although the resonance peaks are strongly dependent on $\tau$, the dip of scattering cross section is almost unchanged, resulting in a relatively robust cloaking. For the dielectric cylinder wrapped by a doped graphene monolayer, the tunability of work frequency of cloaking is dependent on the chemical doping or external voltage gating [28]. In contrast, THz cloaking based on the optically pumped graphene here is manipulated by an all-optical way since the quasi-Fermi energy can be tuned continuously by the intensity of the optical pumped light.

According to equation (9), the tunable cloaking effect here is also related to the properties of dielectric cylinder, as illustrated in figures 3(c) and (d). We calculated the ratio of scattering cross section between the cylinder wrapped by the optically pumped graphene and the corresponding bare one, denoted by $C_{\text{scat}}/C_{\text{ sca}}$. The work frequency of THz cloaking undergoes a red shift as we increase the radius or the permittivity of cylinder. The calculated results indicate that the THz cloaking effect persists as the parameters of dielectric cylinder vary. In addition to the tunable THz cloaking, THz lasing can also be realized in this system. Here, the optically pumped graphene plays two roles, i.e., the plasmonic resonator and gain amplification medium. The contour map of normalized scattering cross section (in log scale) versus the quasi-Fermi energy $\mu_F$ and incident THz frequency is shown in figure 4(a). All the parameters are kept the same as those in figure 2(c) except the quasi-Fermi energy. Remarkably, there exists anomalous scattering peak (THz lasing) at certain range of quasi-Fermi energy. The work frequencies of THz lasing have a blue shift as the quasi-Fermi energy increases. As a result, the THz lasing can be tuned continuously by the quasi-Fermi energy. To be specific, the line-shape of normalized scattering, absorption and extinction cross section for $\mu_F = 0.14$ eV and $\mu_F = 0.19$ eV are given in figures 4(b) and (c), respectively. Note that the negative absorption cross section indicates amplification of light. The scattering enhancement over 4 orders of magnitude is observed at the frequency of 5.17 THz for $\mu_F = 0.19$ eV. Figures 4(d) and (e) show the spatial distribution of magnetic field at the resonance for

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**Figure 3.** Scattering cross section for different quasi-Fermi energy (a) and relaxation time (b) We fix $\tau = 1$ ps in (a) and $\mu_F = 0.16$ eV in (b). The gray dashed lines in (a) and (b) represent the corresponding scattering cross section of a bare dielectric cylinder. We set the permittivity $\varepsilon_1 = 2.1$, and radius $R = 1.5$ $\mu m$ in both (a) and (b). (c) and (d) Ratio of scattering cross section between the coated and bare cylinder with the radius $R$ and permittivity $\varepsilon_1$ varied. The parameters $\mu_F = 0.16$ eV, $\tau = 1$ ps are fixed in (c) and (d).
$\mu_F = 0.14\,\text{eV}$ and $\mu_F = 0.19\,\text{eV}$, respectively. The magnetic field at the vicinity of cylinder has been enhanced with almost 3 orders of magnitude at the resonance frequency for $\mu_F = 0.19\,\text{eV}$.

The effects of dielectric cylinder for the THz lasing are also considered as shown in figures 5(a) and (b). The incident signal THz frequency is fixed to be 5.2 THz without loss of generality. Gigantic enhancement of scattering cross section can be achieved for $\mu_F$ ranging from 0.15 eV $\sim$ 0.2 eV, which means that we can tune the $\mu_F$ (depending on the intensity of pumping power) to realize the THz lasing for different radius $R$ and permittivity $\varepsilon_1$. The quasi-Fermi energy $\mu_F$ for THz lasing increases as the radius of cylinder increases (see figure 5(a)). Thus, the smaller the radius $R$, the lower the power of optical pump required for the THz lasing. Meanwhile, $\mu_F$ should also increase as the permittivity of dielectric cylinder $\varepsilon_1$ increases (see equation (10)). Our calculated results indicate that the THz lasing is available for a wide range of parameters of dielectric cylinders under the proper optical pump.

Furthermore, the work frequency of THz lasing is dependent on the electron scattering rate $\tau$ as shown in figure 5(c). It is found that the work frequency and $\mu_F$ should also increase as $\tau$ increases. The slope of $\mu_F$ and work frequency as a function of $\tau$ for the lasing condition eventually becomes smaller as $\tau$ increases. On the other hand, the work frequency and amplification magnitude of THz lasing is almost unaffected by the temperature of
The reason can be attributed to the fact that $\mu_F = 0.2$ eV we choose is much larger than $k_B T$. This can serve as the room-temperature lasing.

Although the THz lasing in flat graphene with optical pumping has been widely investigated [37, 38], the size of the planar structures are quite large. Here, the cylindrical geometry in our work is sub-wavelength for THz lasing. The work wavelength at lasing ($\sim 5.2$ THz) is almost 20 times larger than the diameter of dielectric cylinder. Moreover, the excitation of SPs in flat graphene is quite difficult due to the mismatch of momentum. In general, periodic gratings [22] or sharp small tips [39] are required to excite graphene SPs. However, the graphene SPs in cylindrical dielectric can be excited easily due to the sub-wavelength feature. Besides, due to the van der Waals forces, several experiments have shown graphene layer can tightly coat the dielectric cylinder [27, 40–43].

4. Conclusion

In summary, we propose an all-optical manipulation of subwavelength THz cloaking and lasing based on a dielectric cylinder wrapped by monolayer graphene. The scattering cross section can be dramatically suppressed or amplified. We show that both the work frequencies of THz cloaking and lasing can be tuned continuously by quasi-Fermi energy of graphene. The properties of THz cloaking and lasing depending on relaxation time of charge carriers and optical parameters of cylinder are also investigated. It is worth mentioning that the pump frequency of IR can be optimized to achieve a high pumping efficiency. Although the radius of dielectric cylinder is deep sub-wavelength for the signal THz frequency, it can be comparable for the infrared wavelength of the pump light. The absorption of graphene at IR frequencies can be greatly enhanced due to the Mie resonance [31], and the corresponding power of IR pump required by the THz lasing can be much lower. Alternatively, as the pumping light is squeezed into the guide mode of dielectric cylinder, it is convenient to integrate with other optical elements. Our proposal here may pave the way to all-optical manipulation of THz lasing and cloaking based on optically pumped graphene.
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