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Graphene-based tunable plasmonic Bragg reflector with a broad bandwidth

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We propose and numerically analyze a plasmonic Bragg reflector formed in a graphene waveguide. The results show that the graphene plasmonic Bragg reflector can produce a broadband stopband that can be tuned over a wide wavelength range by a small change in the Fermi energy level of graphene. By introducing a defect into the Bragg reflector, we can achieve a Fabry–Perot-like microcavity with a quality factor of 50 for the defect resonance mode formed in the stopband. The proposed Bragg reflector could be used as a broadband ultrafast tunable integrated filter and a broadband modulator. In addition, the defect microcavity may find applications in graphene-based resonators. © 2013 Optical Society of America

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Plasmonic nanostructures allow the guiding and manipulation of electromagnetic energy beyond the diffraction limit, enabling applications ranging from integrated photonic circuits [1,2] and photovoltaics [3] to biological sensing [4] and laser beam shaping [5,6]. Metallic plasmonic structures are mainly used for the visible and near-infrared regions because of the optical properties of the noble metals. Recently, graphene has attracted great interest due to its unique electronic and optical properties [7,8]. High carrier mobility of graphene has been utilized in developing ultrafast photonic devices, such as the optical modulator [9] and photodetectors [10,11]. Graphene has also been found as a promising platform for plasmonic applications in the infrared frequency regime. Graphene exhibits a relatively large conductivity that translates into long optical relaxation times ($\tau \sim 10^{-13}$ s), and thus could potentially provide a large plasmon wave propagation distance [12,13]. The carrier density or equivalently Fermi energy level, E_f , relative to the Dirac point of graphene can be adjusted chemically or through bias voltage applied on a field effect transistor (FET) [14] in less than a nanosecond [9]. In addition, graphene plasmons have a strong optical field confinement that has been verified by experiments [15–17]. Graphene plasmon resonances have strong oscillator strengths at room temperature; in comparison, low temperatures (4.2 K) were required for conventional 2D electron gas [18,19]. Much attention has been focused on localized graphene surface plasmon resonance with the incident light from free space, such as nano-ribbons [17,20–22], nano-disks [23], graphene-metal plasmonic antennas [24,25], and graphene metamaterials [26]. A diffractive silicon grating has been demonstrated to excite plasmonic waves in graphene [27]. It is also interesting and indispensable to study the propagation properties of graphene plasmon waves. Guided plasmon waveguiding and hybridization in individual and paired nano-ribbon have been analyzed [13]. Recently, graphene

nano-ribbon bended waveguides and splitters have been investigated [28].

In this work, we first propose a plasmonic Bragg reflector structure formed in graphene waveguides and investigate its performance. We show that periodic stacks of plasmonic graphene–silicon and graphene–air waveguides can be utilized to design effective filtering effects around the Bragg wavelength. The tunability of the filtering stopband by electrostatic and defect cavity modes are also studied. Such graphene Bragg reflectors can be used to build high speed integrated modulators with broadband bandwidth. In addition, we introduce a defect into the Bragg reflector to achieve a defect cavity mode formed in the stopband with a high-Q of 50. Such a defect microcavity may be used to build graphene-based resonators for various applications.

We first investigate the optical properties of surface plasmon waves on a graphene sheet deposited on a silicon and air substrate, as shown in the inset of Fig. 1(b). By solving the dispersion relation of the transverse magnetic (TM) mode in graphene sheet [12], we can obtain the effective refractive index n_{eff} of the surface plasmon waveguide mode on graphene sheet. In the calculations and simulations, graphene sheet is treated as an anisotropic material with thickness $t = 0.5$ nm, where the out of plane permittivity is 2.5 based on the graphite dielectric constant. The optical conductivity of the graphene sheet is calculated with the random-phase approximation (RPA) [29,30] as

$$\sigma(\omega) = \frac{2ie^2k_B T}{\pi\hbar^2(\omega + i\tau^{-1})} \ln \left[2 \cosh \left(\frac{E_f}{2k_B T} \right) \right] + \frac{e^2}{4\hbar} \left\{ \frac{1}{2} + \frac{1}{\pi} \arctan \left(\frac{\hbar\omega - 2E_f}{2k_B T} \right) - \frac{i}{2\pi} \ln \left[\frac{(\hbar\omega + 2E_f)^2}{(\hbar\omega - 2E_f)^2 + (2k_B T)^2} \right] \right\}, \quad (1)$$

Fig. 1. Effective refractive index of a graphene plasmon mode at graphene–silicon waveguide and graphene–air waveguide as a function of optical wavelength for different Fermi energy levels: (a) $\text{Real}(n_{\text{eff}})$, and (b) $\text{Imag}(n_{\text{eff}})$. The depth of the air trench is 20 nm and the carrier mobility used is $\mu = 10000 \text{ cm}^2/(\text{V} \cdot \text{s})$.

where k_B is the Boltzmann constant, T is the temperature, ω is the light frequency, τ is the carrier relaxation time from the impurities in graphene, and $E_f = \hbar V_f (\pi n)^{1/2}$ is the Fermi energy level, where n is the charge carrier concentration, $V_f = 10^6 \text{ m/s}$ is the Fermi velocity, and μ is the carrier mobility in graphene. The first and second terms of Eq. (1) are, respectively, attributed to the intraband transition and interband transition. In our following calculations, the optical conductivity contribution of the interband transition is neglected because the Fermi level is always above half of the photon energy in the simulated midinfrared spectral range, and under this circumstance the intraband transition dominates the conductivity. The in-plane graphene permittivity is characterized by a dielectric function of $\epsilon_{II} = 2.5 + i\sigma(\omega)/(\epsilon_0 \omega t)$. The following calculations are performed below the optical phonon frequency of 1667 cm^{-1} of graphene, so the damping due to the electron–phonon coupling is neglected. The carrier

relaxation time $\tau = \mu E_f / e V_f^2$ is determined by the Fermi energy level and carrier mobility in graphene [20].

Figure 1(a) shows the real part of the effective refractive index of the surface plasmon mode supported by a graphene sheet on silicon and air substrate, for different Fermi energy levels. The effective refractive index at the Fermi energy level $E_f = 0.4 \text{ eV}$ is larger than 100, indicating that the mid-infrared plasmon wavelength is 2 orders of magnitude smaller than its wavelength in free space and the surface plasmon wave is highly localized. As the Fermi energy level increases, the real part of surface plasmon wave effective index decreases, which can be utilized in the tunable plasmonic device in the following sections. In addition, the effective refractive index of the surface plasmon mode on the silicon substrate is larger than that on the air substrate for a given Fermi level energy, which can provide an effective index contrast. The imaginary part of the effective refractive index of the surface plasmon mode at different Fermi energy level is also plotted in Fig. 1(b). One can see that $\text{Imag}(n_{\text{eff}})$ decreases with increasing Fermi energy level, for a given wavelength, indicating that the propagation loss is smaller at higher Fermi energy levels, as the propagation distance is given by $L = 1/[2k_0 \text{Im}(n_{\text{eff}})]$, where $k_0 = 2\pi/\lambda_0$ (wavelength in the vacuum λ_0). As shown in Fig. 1(a), there is a high effective index contrast for the surface plasmon mode between the graphene on silicon and air substrate. Thus, by periodically modulating the effective index along the graphene sheet, which can be realized by alternatively stacking graphene–silicon and graphene–air waveguides, a Bragg reflector will be formed. In addition, a fast tunable device can be expected from the electrostatic tuning properties of the graphene.

Figure 2 shows the schematic of a graphene plasmonic Bragg reflector formed in silicon gratings, for which the Bragg condition is formulated as:

$$L_1 \text{Real}(n_{\text{eff1}}) + L_2 \text{Real}(n_{\text{eff2}}) = m\lambda_b/2, \quad (2)$$

where λ_b is the Bragg wavelength and m is an integer, L_1 and L_2 are, respectively, the lengths of the graphene–silicon and graphene–air plasmonic waveguide, as indicated in Fig. 2. $\text{Real}(n_{\text{eff1}})$ and $\text{Real}(n_{\text{eff2}})$ are the real parts of the refractive indices of the graphene–silicon and graphene–air plasmonic waveguide, respectively. If Eq. (2) is satisfied, surface plasmon wave propagation through the structure at the Bragg wavelength will be prohibited. The finite-difference time-domain (FDTD) simulations are used to calculate the optical response

Fig. 2. Schematic of graphene plasmonic Bragg reflector formed in silicon grating substrate. The depth of air trench is 20 nm.

Fig. 3. (a) Simulated transmission spectra of graphene plasmonic reflector with different Fermi energy levels when $\mu = 10000 \text{ cm}^2/(\text{V} \cdot \text{s})$; (b) Bragg wavelength as a function of Fermi energy level. The squares are simulated results from FDTD, and the red line is calculated from Eq. (2). The lengths of the graphene-silicon, graphene-air waveguide are $L_1 = 80 \text{ nm}$, $L_2 = 40 \text{ nm}$, respectively. The period of Bragg cell N is 10.

of the graphene Bragg reflector. The perfectly matched layer absorbing boundary condition is set at all boundaries of the simulated domain.

Figure 3(a) shows the simulated transmission spectra of the graphene plasmonic reflector with the Fermi energy levels $E_f = 0.6 \text{ eV}$, 0.7 eV , and 0.8 eV , and period of Bragg cell $N = 10$. One can see that there are two wide stopbands with near-zero transmissions around the Bragg wavelength of $9.1 \mu\text{m}$ (corresponding to $m = 2$) and wavelength of $7.3 \mu\text{m}$ (corresponding to $m = 3$) for $E_f = 0.6 \text{ eV}$ which shows good filtering characteristics. The bandwidths of the stopband (defined as the difference between the two wavelengths at each of the transmittance is equal to 1%) are, respectively, 368 and 310 nm for the Bragg wavelength of 9.1 and $7.3 \mu\text{m}$. Broadband and faster electrostatic tunability are the most intriguing properties of graphene materials. Simulated spectra in Fig. 3(a) shows that the stopband shifts to a shorter wavelength with an increase of the Fermi energy level. Figure 3(b) shows the central wavelength of Bragg reflector as a function of Fermi energy levels, which confirms the broad tuning range with a small change of Fermi level. The FDTD simulations results agree very well with the theoretical curve obtained from Eq. (2).

Figure 4 shows the transmission spectra of the graphene Bragg reflector for different period numbers.

Fig. 4. Transmission spectra of a graphene Bragg reflector consisting of 6, 8 and 10 periods with Fermi energy level $E_f = 0.8 \text{ eV}$, and reflection spectrum for periods of 10. Other structure parameters are the same as in Fig. 3.

In principle, Bragg scattering exists at any number of Bragg periods. In the case of the graphene Bragg reflector, it is found that the minimum N required to realize transmission less than 0.4% in the stopband is $N = 10$. The spectrum displays some sidelobes out of the stopband which, we believe, are due to light scattering at the abruptly disappearing boundary at the end of the Bragg gratings. As shown in Fig. 4, the increase of period number gives rise to higher propagation losses. The reflectance at the Bragg wavelength of $7.9 \mu\text{m}$ is as high as 96.4%, while the transmittance is below 0.4%. In addition, we calculate the transmittances of the graphene Bragg reflector at different carrier relaxation times and find that the transmittances at the wavelength of $8.4 \mu\text{m}$ are 0.4%, 25%, 49% for $\tau = 1.6 \times 10^{-13} \text{ s}$, $4 \times 10^{-13} \text{ s}$, and $8 \times 10^{-13} \text{ s}$, which are corresponding to the carrier mobilities of $2000 \text{ cm}^2/(\text{V} \cdot \text{s})$, $5000 \text{ cm}^2/(\text{V} \cdot \text{s})$, and $10000 \text{ cm}^2/(\text{V} \cdot \text{s})$ of graphene with periods $N = 10$, Fermi energy level $E_f = 0.8 \text{ eV}$. One can conclude that, with a large relaxation time, the graphene Bragg reflector has a lower propagation loss.

To further investigate the properties of the graphene plasmonic Bragg reflector, we introduce a defect into the graphene Bragg plasmonic structure by decreasing the length L_d of the central graphene-silicon waveguide. Figure 5(a) shows the transmission spectrum when the 4th graphene-silicon waveguide length is reduced to 50 nm. The cavity length Q -factor is defined as $Q = \lambda_0/\Delta\lambda$, where λ_0 and $\Delta\lambda$ are, respectively, the central resonance wavelength and full width at half-maximum of the defect mode. It describes the ratio of the energy stored in the microcavity at resonance to the energy escape from the cavity per cycle of oscillation. The Q -factor for the cavity in Fig. 5(a) is found to be about 50, which is higher than the nano-ribbon resonance cavity [17] in the literature. One can see that the stopband of the defect cavity is a little larger than the Bragg reflector, which is attributed to the destructive interference inside of the Fabry-Perot-like microcavity. As shown in Fig. 5(b), the peak wavelength of the defect mode can be adjusted by changing the length of the cavity. In addition, the peak wavelength can also be easily shifted by electrically tuning the graphene carrier density.

Figure 6 shows the field profiles of the graphene surface plasmon propagation through the Bragg reflector

Fig. 5. (a) Transmission spectra of a Bragg reflector with period number $N = 8$ and the microcavity formed by introducing a defect length $L_d = 50 \text{ nm}$; (b) peak wavelength as a function of defect length. The Fermi energy level $E_f = 0.8 \text{ eV}$, and the lengths of the graphene-silicon, graphene-air waveguide are $L_1 = 80 \text{ nm}$, $L_2 = 40 \text{ nm}$, respectively.

Fig. 6. (a) and (b) Simulated electric field intensity profiles of the plasmonic Bragg reflectors at the incident wavelength at 10.5 and 9.1 μm , and the structure parameters are the same as Fig. 3(a) with Fermi energy level of 0.6 eV; (c) corresponds to the defect resonance mode in Fig. 5(a) at 7.6 μm .

and the defect microcavity. Figure 6(b) shows that the incident wave is reflected at the wavelength of 9.1 μm , while it transmits through the structure at the wavelength of 10.5 μm , which is out of the stopband [Fig. 6(a)]. As displayed in Fig. 6(c), the surface plasmon wave at the wavelength of 7.6 μm couples into the microcavity structure, which agrees well with the spectrum shown in Fig. 5(a).

In conclusion, we have proposed and numerically analyzed plasmonic Bragg reflectors formed in graphene waveguides. The plasmonic Bragg reflector can achieve a broadband fast-tunable stopband. This structure can also be used as a modulator for a Bragg reflector with structure parameters in Fig. 3; a modulation depth of 21.4 dB can be achieved at the wavelength of 8.34 μm when the Fermi level is increased from 0.7 to 0.8 eV, which is the appealing property of the device compared with metallic plasmonic Bragg structures [31,32]. By introducing a defect in the Bragg reflector, we can obtain a defect resonance mode with a high-Q factor. Those proposed structures could find applications in building active integrated photonic circuits.

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