Research article

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Wideband tunable perfect absorption of graphene plasmons via attenuated total reflection in Otto prism configuration

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Abstract: A strategy is proposed to achieve wideband tunable perfect plasmonic absorption in graphene nanoribbons by employing attenuated total reflection (ATR) in Otto prism configuration. In this configuration, the Otto prism with a deep-subwavelength dielectric spacer is used to generate tunneling evanescent waves to excite localized plasmons in graphene nanoribbons. The influence of the configuration parameters on the absorption spectra of graphene plasmons is studied systematically, and the key finding is that perfect absorption can be achieved by actively controlling the incident angle of light under ATR conditions, which provides an effective degree of freedom to tune the absorption properties of graphene plasmons. Based on this result, it is further demonstrated that by simultaneously tuning the incident angle and the graphene Fermi energy, the tunable absorption waveband can be significantly enlarged, which is about 3 times wider than the conventional cavity-enhanced configuration. Our

proposed strategy to achieve wideband, tunable graphene plasmons could be useful in various infrared plasmonic devices.

Keywords: graphene; plasmons; perfect absorption; attenuated total reflection; wide waveband.

1 Introduction

Graphene, the truly two-dimensional crystal consisting of carbon atoms in a honeycomb lattice, has received extensive research interest due to its excellent optoelectronic properties and potential applications [1, 2]. Its unique electronic features, characterized by massless Dirac fermions and unconventional transport properties, render it to possess the most distinct electromagnetic confinement behavior: the capability of supporting intrinsic plasmon modes with extremely tight confinement, relatively long lifetime, and very low losses at infrared and terahertz frequencies [3–11]. Especially, such plasmonic responses can be actively modulated under electrostatic gate control [12–17], a feature that is not available in traditional metallic surface plasmons. These unique properties make graphene a promising one-atom-thick platform for the realization of highly integrated active plasmonic devices, such as infrared biosensors [18–22], plasmonic modulators [23], plasmonic infrared photodetectors [24–26], plasmonic thermal emitters [27], and so forth. Despite these exciting potentials, the relatively low excitation efficiency of graphene plasmons is a major obstacle for graphene-based optoelectronics, due to the inherent one-atom thinness of graphene as well as the large wave vector mismatch between graphene plasmons and free-space photons [28].

To address this issue, strategies such as increasing graphene's doping concentration [16], stacking graphene multilayers [29–31], and utilizing external photonic structures [32–38] have been utilized to realize the efficient coupling of free-space photons to graphene plasmons. Among these strategies, utilizing external photonic structures is recognized as the most efficient

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way to achieve stronger graphene-light interaction. It has been theoretically predicted that perfect absorption (100% efficiency) is possible in graphene nanostructures positioned above a Fabry-Pérot cavity comprised of an opaque reflector and a lossless dielectric spacer [32, 33, 39]. The strategy is to overcome the inefficient coupling between free-space photons and highly confined graphene plasmons through multiple reflections. Specifically, at the Fabry-Pérot cavity resonant frequencies, the metal reflector can create a standing wave between the incident and reflected light, which maximizes the electric field overlapping with graphene. Once the graphene plasmonic frequency is tuned to coincide with such frequencies, a significant improvement of graphene plasmonic absorption can occur [28]. It is worth pointing out that perfect absorption can be realized in graphene with low carrier mobility by further integrating graphene plasmonic Fabry-Pérot resonators with sub-wavelength metallic slit arrays, which has been recently realized experimentally [37]. However, perfect absorption can occur only at a specific frequency, and once the graphene plasmonic frequency is tuned away from the cavity resonant frequency, the absorption significantly decreases [40]. In other words, the tunable bandwidth of such resonant-cavity-enhanced graphene plasmons is limited, which restricts its practical use in various applications that require a widely tunable waveband. Recently, wideband absorbers were realized at terahertz frequencies based on strategies such as stacked multilayered graphene nanostructures [41], graphene plasmonic structures with gradient widths [42, 43], and net-shaped periodic sinusoidally patterned graphene [44]. They are promising to enlarge the bandwidth but require subtle device fabrication techniques.

In this article, we propose a nonresonant strategy to realize perfect absorption over a wide waveband via attenuated total reflection (ATR) by placing graphene nanoribbons in an Otto prism configuration with a deep-subwavelength space layer. In this configuration, graphene nanoribbons are used to provide a sufficient wave vector for the direct excitation of graphene plasmons even without the prism. The prism is employed to block the transmitted “channel” so that there exists only the incident/reflected “channel” for the light to couple with graphene nanoribbons to excite plasmons, resulting in high absorption of the incident infrared light. It is demonstrated that by simultaneously adjusting the incident angle and the Fermi energy of graphene, the tunable wavelength region with high plasmonic absorption can be much wider than that of the conventional Fabry-Pérot cavity-based configuration.

2 Structure and principle

The schematic of the Otto configuration used to realize perfect absorption of graphene plasmons via ATR is sketched in Figure 1A. From top to bottom are shown graphene nanoribbons, the dielectric spacer layer, and the prism. It is assumed that the refractive index of the spacer is smaller than that of the prism, i.e. $n_s < n_p$. Under this assumption, the physical processes of graphene plasmon excitations can be categorized into three regions depending on the incident angle of the light.

Region I is illustrated in Figure 1C, where the incident angle θ is smaller than the critical angle of ATR at the spacer/air interface θ_{ca} (i.e. $\theta < \theta_{ca} \equiv \arcsin(n_a/n_p)$, with n_a , n_p being the refractive index of air and prism, respectively). In this region, ATR cannot occur at either the spacer/air interface or the prism/spacer interface, and hence part of the incident light is reflected, and the other part is transmitted, which would further couple with graphene nanoribbons to excite plasmons of the ribbons. From the perspective of the temporal coupled mode (TCM) theory, only in the case of one coupling “channel” can perfect absorption (100% absorption) be achieved [45] (details can be seen in Supplementary Information). Region I has two “channels” to couple with graphene plasmons: i.e. the incident/reflected “channel” and the transmitted “channel”, and therefore the absorption of graphene plasmons cannot reach 100%, no matter how we optimize the configuration.

Region II corresponds to the situation where the incident angle θ fulfills the relation $\theta_{ca} < \theta \leq \theta_{cs} \equiv \arcsin(n_s/n_p)$, with θ_{cs} being the critical angle of ATR at the prism/spacer interface and n_s the refractive index of spacer, as shown in Figure 1D. In this case, ATR can occur at the spacer/air interface, and thus there would be an evanescent wave at the interface, which will couple with the graphene ribbons to excite the plasmons. According to TCM theory, now only the incident/reflected coupling “channel” exists, and one can achieve 100% absorption of the incident light by optimizing the configuration.

Region III is illustrated in Figure 1E, where the incident angle θ is larger than θ_{cs} so that ATR at the prism/spacer interface can occur. This case is similar to Region II, except that the coupling between the evanescent wave at the prism/spacer interface and graphene plasmons becomes weaker, because the magnitude of the evanescent wave exponentially decreases away from the interface. However, the absorption of graphene plasmons can still reach 100% according to TCM theory because there exists only the incident/reflected “channel”, which will be discussed in detail in Section 3.

Figure 1: The physical processes of graphene plasmon excitations using Otto configuration.

(A) Schematic of the structure consisting of patterned graphene nanoribbons on a prism separated by a spacer layer. (B) Electric field $|E_x|$ of the mode at the resonant wavelength. Absorption, transmission, and reflection spectra of the structure for the cases (C) I, (D) II, and (E) III, when the incident angle is 10° , 40° , and 70° , respectively. The inset shows the cross-section of the structure.

To intuitively understand the above discussions, we simulated graphene plasmon excitation processes by finite element method using Comsol Multiphysics. The surface conductivity $\sigma(\omega)$ of graphene is modeled by a Drude-like formula in the considered infrared region as [46, 47]

$$\sigma(\omega) = \frac{e^2 E_f}{\pi \hbar^2} \frac{i}{\omega + i\tau^{-1}}. \quad (1)$$

In Eq. (1), e is the elementary charge, E_f is the Fermi energy, and $\hbar$ is the reduced Planck constant.

The electron relaxation time τ is defined as $\tau = \mu_c E_f / e v_f^2$, where $v_f = c/300$ is the Fermi velocity, with c being the speed of light in vacuum, and μ_c is the carrier mobility in graphene. In the simulations, the Bloch periodic boundary condition is imposed in the x -direction, while the perfectly matched layer (PML) is applied in the z -direction to achieve the absorbing boundary conditions at two ends of the computational space. Nonuniform meshes are used in the simulation regions, where the mesh size gradually increases outside the graphene. A TM-polarized light with input electric field intensity of 1 V m^{-1}

in a period is incident upon the structure with incident angle θ . The period and width of the graphene nanoribbon are fixed as $\Lambda=100$ nm and $W=50$ nm. A moderate graphene carrier mobility of $\mu_c=10,000$ cm² V⁻¹ s⁻¹ and Fermi energy of $E_f=0.4$ eV are used in the simulation [48]. The refractive index of the spacer and the prism are $n_s=1.5$ and $n_p=2$, respectively. The thickness of the spacer is $t=50$ nm.

The simulation results are presented in the lower panels of Figure 1C–E. When the incident angle $\theta=10^\circ$ (i.e. in Region I), an absorption peak with magnitude about 0.65 at 7.25 μm is observed in the absorption spectrum, as shown in Figure 1C. The electric field mode profile $|E_x|$ at this peak is plotted in Figure 1B, from which we can see that a highly confined graphene plasmonic mode is excited. The electromagnetic energy of the mode is dissipated in the graphene nanoribbon due to Ohmic losses of free electron oscillations, yielding the peak in the absorption spectrum. Another feature in Figure 1C is that the transmission and reflection spectra also change drastically at the resonant wavelength of 7.25 μm . As the incident angle θ increases to 40° (i.e. in Region II as shown in Figure 1D), ATR occurs at the spacer/air interface and the transmission is totally blocked. In this situation, perfect absorption is observed at the resonant wavelength of 7.25 μm . These results are in accordance with TCM theory, which shows that Region II can reach perfect absorption because there is only one coupling “channel”. As θ further increases to 70° (Region III, Figure 1E), ATR occurs at the prism/spacer interface, and the incident light first tunnels into the evanescent wave and then gets coupled to the graphene plasmonic wave. Note that the absorption in this case is about 90%, which is slightly lower than in Region II, due to the weaker coupling between the ATR evanescent wave and graphene plasmons, but perfect absorption can still be achieved by optimizing the parameters of the configuration, as we discuss in the next section.

The resonant wavelength in the simulated spectra can be predicted by the dispersion relation of the localized plasmons of graphene ribbons, given by

$$\lambda_{\text{GPs}} = \frac{\pi\hbar c}{e} \sqrt{\frac{2 \times \xi \epsilon_0 (\epsilon_a + \epsilon_{\text{eff}}) \Lambda}{E_f}}. \quad (2)$$

Here, λ_{GPs} is the resonant wavelength of graphene plasmons, ϵ_0 is the vacuum permittivity, ϵ_a is the relative permittivity of air, ξ is a dimensionless fitting constant, and ϵ_{eff} is the effective permittivity of the spacer and prism when they are taken as an effective uniform medium, which can be approximated as [49, 50]

$$\epsilon_{\text{eff}} \equiv \epsilon_s \frac{(\epsilon_p + \epsilon_s) \cdot \exp(2qt) + (\epsilon_p - \epsilon_s)}{(\epsilon_p + \epsilon_s) \cdot \exp(2qt) - (\epsilon_p - \epsilon_s)}. \quad (3)$$

In Eq. (2), ϵ_p and ϵ_s are the relative permittivity of the prism and the dielectric spacer, respectively, and q is the wave vector of graphene plasmon waves. Based on Eqs. (2) and (3), we can calculate that the resonant wavelength is 7.25 μm if the fitting parameter $\xi=1.51$, in accordance with the simulation results.

3 Results and discussion

In this section, we study the dependence of the absorption spectra on the incident angle, spacer thickness, and Fermi energy of graphene. Based on these results, we further investigate how to achieve wideband tunable perfect absorption of graphene plasmons in such configurations.

3.1 Tunable absorption by Incident angle

At first, the evolution of absorption, transmission, and reflection spectra as a function of the incident angle is investigated. The simulated spectra with varying incident angle are shown in Figure 2A–C respectively. It can be seen that the resonant wavelength of graphene plasmons is nearly independent of the incident angle. This is because the resonant wavelength of localized graphene plasmons is mainly determined by the geometric parameters of graphene nanoribbons and the Fermi energy. Meanwhile, a clear dividing line can be observed at the incident angle of 30° in the three figures, corresponding to the critical angle of ATR occurring at the spacer/air interface (further proof can be seen in Figures S2 and S4). When the incident angle is larger than this critical angle, total reflection occurs and perfect absorption is realized at the resonant wavelength of 7.25 μm .

To get a clearer view of the above result, the resonant absorption, transmission, and reflection at 7.25 μm with varying incident angles are extracted from the three figures and shown in Figure 2D. Three distinct regions corresponding to the three cases discussed previously can be observed. For the Region I ($0^\circ < \theta \leq 30^\circ$), transmitted and reflected light can leak out of the structure, resulting in a relatively small resonant absorption. For the Region II ($30^\circ < \theta \leq 48.6^\circ$), the transmission is blocked and nearly perfect absorption is observed within a wide incidence angle range of 35° – 49° . For Region III ($\theta > 48.6^\circ$), the

Figure 2: Graphene plasmonic absorptions versus the incident angle.

Absorption (A), transmission (B), and reflection (C) mapping as a function of the incident angle and wavelength when $n_s = 1.5$, $n_p = 2$. (D) Extracted resonant absorption, transmission, and reflection with varying incident angles.

evanescent field intensity decreases as the incident angle increases and exponentially decays away from the interface, which degrades the excitation of graphene plasmons. Hence, perfect absorption is easier to achieve in Region II since graphene nanoribbons are located at the ATR interface where the intensity of the evanescent field is stronger than that of Region III, but in Region III perfect absorption can still be observed in a small angle range (i.e. 48.6° – 55°). It is worth noting that a sharp decrease of absorption can be observed at the first critical angle (30°). This phenomenon is caused by the change of the electric field intensity at the spacer/air interface (details can be seen in Figure S3).

Therefore, it becomes clear that the absorption, transmission, and reflection spectra of graphene plasmons in our configuration have a strong dependence on the incident angle of the infrared light. By varying the incident angle, the excitations of graphene plasmons can be categorized into three regions with different absorption properties, representing a useful degree of freedom to tune the properties of graphene plasmons.

3.2 Effect of spacer thickness

Next we investigate the effect of the spacer thickness t on the absorption properties of the configuration. The absorption spectra with varying spacer thicknesses t for a fixed incident angle of $\theta = 40^\circ$ (Region II) are shown in Figure 3A. Clearly, the resonant wavelength of the configuration blue-shifts first and then remains unchanged with increasing t . To be more specific, we extracted the resonant wavelengths and absorptions from Figure 3A and plotted them as a function of the spacer thickness in Figure 3B. We can see that the resonant wavelength shifts from 8.9 to 7.3 μm as t increases from 0 to 30 nm, but remains nearly unchanged as t increases further. The blue shift of the resonant wavelength can be explained by Eqs. (2) and (3), that the increase of t results in the decrease of the effective permittivity ϵ_{eff} below that of graphene and hence results in the blue shift of the resonant wavelength. Also, this change in effective permittivity significantly affects the mode pattern of graphene plasmons, as shown in the inset of Figure 3B, which results in the

Figure 3: Effect of spacer thickness on the plasmonic absorption.

(A) Absorption mapping as a function of the spacer thickness and incident wavelength. (B) Extracted resonant wavelength and absorption with varying spacer thicknesses. The green line is the analytically calculated results from Eq. (2). The inset shows the mode pattern of graphene plasmons. (C) Absorption mapping as a function of the incident angle and wavelength when $t = 0.5, 1$, and 2 μm . (D) Extracted resonant absorption with varying incident angles when $t = 0.5, 1$, and 2 μm . The refractive index of the prism is $n_p = 2$.

redistribution of the electric field, especially for a small spacer thickness. However, as t further increases, the perturbation of the prism on the electric field becomes weak, and the effective permittivity ϵ_{eff} has a negligible dependence on t because in Eq. (3) $(\epsilon_p + \epsilon_s) \exp(2qt) \gg (\epsilon_p - \epsilon_s)$ for large t , and thus Eq. (3) becomes $\epsilon_{\text{eff}} \approx \epsilon_s$. Therefore, the resonant wavelength remains unchanged for large spacer thicknesses. The simulated resonant wavelength can be well reproduced by Eqs. (2) and (3), as shown in Figure 3B. In this case, the resonant wavelength is also independent of the prism's refractive index (Figure S4), which allows perfect absorption with lower graphene mobility by using a prism of smaller refractive index (Figure S5). Another observation from Figure 3B is that perfect absorption can be achieved independent of the spacer thickness, as shown in the blue hollow line in the figure. The insensitivity of resonant absorption on the spacer thickness of such a configuration makes it easier to fabricate in practice.

Further simulation results indicate that our configuration can work with high absorption in a broader range

of incident angles for smaller spacer thicknesses. The absorption spectra with varying incident angles when $t = 0.5, 1$ and 2 μm are compared in Figure 3C. Obviously, strong resonant absorption can still be achieved for all the cases when the incident angle is slightly larger than the critical angle θ_{ca} of the spacer/air interface. However, as the incident angle increases, the resonant absorption significantly decreases, which is quite obvious especially for the case when $t = 2$ μm . Quantitatively, the resonant absorption is extracted from Figure 3C and plotted as a function of incident angle in Figure 3D. Perfect absorption can be achieved in the incident angle range 34° – 40° for all the cases. As the incident angle increases to 60° , the resonant absorption is still more than 90% for $t = 0.5$ μm . In contrast, the resonant absorption rapidly drops to only 20% for $t = 2$ μm . Such difference is reasonable, because the intensity of the tunneling evanescent wave decreases as the incident angle increases, which makes it unable to effectively couple with graphene plasmons for larger spacer thicknesses. Therefore, a smaller spacer thickness

is favorable, since it facilitates the coupling of the incident light with graphene plasmons over a wide incident angle range.

3.3 Wideband tunable perfect absorption

Finally, we investigate the band tunability of the configuration by varying the Fermi energy of graphene, which can be realized via chemical or electrostatic doping in practice. The simulation parameters are set the same as in Figure 2, except that the period and the width of the graphene nanoribbons are $\Lambda=200$ nm and $W=100$ nm. The absorption spectra with varying Fermi energy values for a fixed incident angle of $\theta=40^\circ$ are presented in Figure 4A. We can see that the resonant wavelength blue-shifts from 14.6 to 6.5 μm as the Fermi energy increases from 0.2 to 1 eV, which shows good agreement with the analytical dispersion relations of Eq. (2) (Figure 4B). This result indicates that the resonant wavelength can be tuned dynamically

by the Fermi energy without changing the system configuration, which is of great importance in practical applications. Besides the evolution of the resonant wavelength, the resonant absorption also varies with Fermi energy. To have a clearer view of this, we extracted the absorption spectra with varying Fermi energy values from 0.2 to 1 eV in intervals of 0.1 eV and plotted them in Figure 4C. We can see that as the Fermi energy increases, the resonant absorption peak first increases, reaches the maximum value of 100% at 10.1 μm , and then decreases. The resonant peaks extracted from Figure 4A show that high absorption with values larger than 90% can be achieved over a wide wavelength range (8.7–12.1 μm), which can be useful in many practical applications.

We also compared the band tunable properties of our Otto prism configuration with those of the cavity-enhanced configuration, as shown in the inset of Figure 4D, where a Fabry-Pérot cavity is formed between the graphene nanoribbon and the metallic reflective layer to enhance the graphene plasmon absorption. To make a

Figure 4: Tuning the plasmonic absorption by varying the Fermi energy of graphene.

(A) Absorption mapping as a function of the Fermi energy and incident wavelength. (B) Comparison between numerical (hollow spheres) and analytical (black line) resonant wavelength with varying Fermi energy. (C) Absorption spectra with varying Fermi energy when $\theta=40^\circ$. The refractive index of prism is $n_p=2$, and the graphene carrier mobility is $\mu_c=10,000$ $\text{cm}^2 \text{V}^{-1} \text{s}^{-1}$. (D) Absorption spectra with varying Fermi energy for the cavity-enhanced configuration when the spacer thickness is set as $T=0.9\mu\text{m}$. The inset shows the schematic of the cavity-enhanced configuration.

fair comparison, the parameters of graphene nanoribbons (i.e. the width, period, and mobility) in the cavity configuration are set the same as with our prism configuration. The spacer thickness between the graphene nanoribbon and reflective layer in the cavity configuration is set to be $0.9\ \mu\text{m}$ to ensure that perfect absorption occurs at $10.1\ \mu\text{m}$, the same as with the prism configuration. Then, we simulated the absorption spectra of the cavity-enhanced configuration. The resultant absorption spectra with varying Fermi energy is presented in Figure 4D. It shows that high absorption with values greater than 90% can only be maintained in the range $9.4\text{--}11.5\ \mu\text{m}$, which is considerably narrower than that of our prism configuration. Such a result can be understood from the fact that only when the graphene plasmonic wavelength is tuned to coincide with the cavity resonant wavelength can the absorption be significantly enhanced. Once the graphene plasmonic wavelength is tuned away from the cavity resonant

wavelength, absorption would decrease rapidly. Therefore, our prism configuration shows obvious advantages over the cavity-enhanced configuration in terms of the tunable wavelength region with high absorption.

In fact, the tunable wavelength region with high absorption of our prism configuration can be further enlarged, because we have another degree of freedom to tune the absorption, namely the incident angle of light (as discussed previously). In what follows, we show that the tunable wavelength region can be significantly enlarged by simultaneously tuning the incident angle and the Fermi energy of graphene. To give an intuitive illustration, the absorption as a function of the incident angle and wavelength when the Fermi energy increases from 0.2 to 1 eV is mapped in Figure 5A, and the corresponding extracted resonant absorptions as a function of Fermi energy at some representative ATR angles are plotted in Figure 5B. It is observed that the resonant absorption first increases,

Figure 5: Wideband tunable perfect absorption by varying incident angle and graphene Fermi energy simultaneously.

(A) Absorption mapping as a function of the incident angle and wavelength when $E_f = 0.2, 0.4, 0.6, 0.8$, and 1 eV. (B) Extracted resonant absorption as a function of Fermi energy with varying incident angles. (C) Absorption spectra with varying Fermi energy under optimized angles in the critical coupling regime.

reaches perfect absorption, and then decreases as the Fermi energy increases. This corresponds to the evolution of the system from an “under-coupling” state, through a “critical coupling” state, and finally to an “over coupling” state with the increase of Fermi energy. The physics behind such phenomenon is that, as the Fermi energy increases (i.e. graphene electron density increases), the coupling strength between graphene plasmons and the incident light gets larger, resulting in the increase of the coupling rate. On the other hand, the loss rate of graphene plasmons will become smaller because the electron screening effect that contributes to the reduction of plasmon loss will be stronger, which is captured in our simulation by the carrier lifetime $\tau = \mu_c E_f / e v_f^2$ in the Drude-like mode of graphene conductivity [Eq. (1)]. Hence, it is the changes of the coupling rate and loss rate with the variation of Fermi energy that drive the system into different coupling states. By simultaneously controlling the Fermi energy and the incident angle, we can always guarantee that the system is in the “critical coupling” state so that the absorption is high. For example, as the Fermi energy increases from 0.4 to 1 eV, we can choose the critical angle to ensure that the system is in the “critical coupling” state to achieve wideband, tunable, perfect absorption for all Fermi energy values, as shown in Figure 5C. Hence, the range with absorptions >90% is significantly enlarged, which is now 6–12 μm as the Fermi energy increases from 0.3 to 1 eV, which is about 1.77 times wider than for the same prism configuration without tuning the incident angle, and 2.86 times wider than for the cavity-enhanced configuration.

Finally, we would like to point out that our proposed Otto prism configuration is essentially different from the configuration where graphene plasmons are excited by a prism, as discussed in Refs. [51–54]. For the latter configuration, a high-refractive-index prism is used to provide a sufficient wave vector to excite plasmons in a continuous graphene sheet, and thus the resonant wavelength is sensitively dependent on the incident angle. In contrast, for our proposed configuration, plasmons of graphene ribbons can be excited directly even without the prism. The role of the prism is to block the transmitted “channel” so that there exists only the incident/reflected “channel” for the light to couple with graphene nanoribbons to excite plasmons. Therefore, the resonant wavelength is independent of the incident angle, which provides us an extra degree to tune the resonant condition to achieve perfect absorption while keeping the resonant wavelength unchanged.

The proposed ATR configuration with impressive features points to many promising applications of graphene plasmons. For example, it could be employed for surface-enhanced infrared absorption spectroscopy for in situ

aqueous-phase biosensing. The background absorption of bulk water molecules can be suppressed using the ATR configuration, which is not possible in the conventional transmission-type configurations [18–22]. This is extremely important for bioanalysis, since biomolecules are active only in aqueous media and many chemical reactions associated with life can take place only in aqueous solutions.

4 Conclusions

In summary, we proposed a strategy to achieve perfect absorption of graphene plasmons in a wide waveband via ATR using an Otto prism configuration. We demonstrated that perfect absorption can be achieved as the incident angle is tuned to the ATR region (i.e. Regions II and III). The influence of spacer thickness on the absorption behavior of graphene plasmons was also studied in detail. It was found that a relatively thin spacer layer is preferable for assisting the coupling between evanescent waves and graphene plasmons. We further demonstrated that by simultaneously controlling the incident angle of light and the Fermi energy of graphene, the tunable wavelength region with high absorption can be significantly enlarged, which is about 3 times wider than that of the conventional Fabry-Pérot cavity-enhanced configuration. Our study provides a promising strategy to achieve the wideband, tunable, perfect absorption of graphene plasmons, and could be beneficial to various infrared graphene plasmonic applications.

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