

1. Introduction

With increasing demand for the implementation of novel optoelectronic applications based on the optical properties of semiconductors, the development of nanoscale semiconductors exhibit shape-, size-, and composition-dependent electronic and optical properties. Unfortunately, when the dimensions of semiconductors are reduced, light-matter interactions will be weakened accordingly, which is adverse to optoelectronic applications. In order to improve the light-matter interactions, numerous approaches have been presented, such as doping metal into semiconductor nanostructures [1, 2], synthesizing heteronanostructures [3–5], and coupling semiconductor nanomaterials with metal nanoparticles (MNPs) [6–9] to form hybrid nanostructures. The hybrid nanostructures exhibiting unique properties in comparison to their individual counterparts have been investigated both theoretically and experimentally [10–13].

Benefiting from advances in modern nanoscience and nanotechnology, the construction of various nanostructures such as semiconductor quantum dots (SQDs) and MNPs for their applications in photonics and optoelectronics have attracted tremendous attention. On one hand, SQDs, as a simple stationary atom with the tunability of optical properties [14] exhibiting abundant physical phenomena of quantum-confined systems [15, 16], are attractive for they present
optical responses in a widely tuned spectral range and easily couple to other nanomaterials [17], which paves the way for numerous potential applications [18]. On the other hand, MNPs sustain collective electronic excitations (i.e. localized surface plasmons) that can strongly enhance electromagnetic fields [19–21], thus enhancing their ability for focusing optical fields. Currently, a significant amount of hybrid nanostructures focusing on hybrid artificial molecule composed of SQD-MNPs systems have attracted significant interest [7, 19, 22, 23], which profoundly modifies light-matter interactions with potential tunability and control. The SQD proximity to plasmonic nanostructures will induce significant alteration of the electromagnetic field felt by the SQD due to the interaction between the excitons (discrete quantum confined electronic states in the SQD) and the surface plasmons dielectric-confined electromagnetic modes in the MNP [24–26].

Due to the exciton-plasmon interaction, several interesting phenomena, such as energy transfer [27], intrinsic optical bistability [28], and manipulation of population inversion [29], have been explored in the hybrid complex SQD-MNP systems. The hybrid system also provides an intuitive picture for highly sensitive detections and eventually leads to innovative new devices, such as applications of DNA sensors [30], laser systems without cavities [31], and the manipulation of heat generation in MNPs [6].

For theoretically interpreting the exciton-plasmon interaction in the hybrid SQD-MNP systems, there are two representative theory treatments at present. The first one is a semiclassical treatment, where the exciton was described in the quantum framework, while the description of the plasmon was within the classical electromagnetic dynamics [9, 32]. Based on the semiclassical description, many properties and phenomena of the hybrid system have been revealed, such as Fano effect [33, 34], nonlinear Fano effect [6], plasmonic electromagnetically induced transparency [35, 36], slow light [37], and superluminal light propagation [38]. Although the semiclassical treatment is generally believed to be correct in the weak-field limit, the nonlinear behavior such as optical bistability will be washed out by quantum fluctuations in the strong-field regime. Then the second one of a quantum treatment is called out [7, 23, 39–42]. In this treatment both the exciton and the electromagnetic field of the MNP are quantized, enabling a more precise analysis of their mutual interactions, and quantum description of the hybrid systems paves the way for potential applications in quantum information processing [43]. In addition, Fano effect from the quantum description differs both qualitatively and quantitatively from the semiclassical description, especially in the strong field regime, which was also demonstrated [44]. Therefore, quantum description for the exciton-plasmon interaction can reveal more novel optical properties that may be applied in optical processing devices in the future.

Although the phenomena of electromagnetically induced transparency (EIT) [35, 36, 45], slow light [37], and superluminal light propagation [38] are investigated in the hybrid SQD-MNPs systems, the treatment is still in the semiclassical regime. Fano effect [33, 34] is demonstrated in the quantum description, however Fano resonance induced the coherent optical propagation in the quantum treatment, such as the fast light and slow light effects, even the switch of fast-to-slow light has never attracted enough academic attention. In this work, we consider a hybrid artificial molecule of the SQD-MNP system, and investigate the coherent optical propagation using the principles of cavity quantum electrodynamics (C-QED) [7, 23, 46]. In this context, the MNP in the hybrid SQD-MNP system plays the role of the cavity in C-QED [47]. We first demonstrate that the probe absorption spectra of the QD can display the switch from EIT to Fano resonances induced by the exciton-plasmon interaction under different detuning regimes, which can be explained by the interference effect in terms of the dressed states. The Fano resonances can be effectively tuned and the probe absorption spectra can display a series of asymmetric Fano line shapes under different parameter regimes including exciton-plasmon coupling strengths $g$, the exciton-pump field detuning $\Delta_\text{ex}$, and the surface-plasmon polaritons (SPP)-pump field detuning $\Delta_\text{SPP}$. Secondly, we investigate the coherent optical propagation properties including the effects of fast and slow light by numerically calculating the group delay of the probe field around the transparency window accompanied by the steep phase dispersion. The results even indicate that a tunable and controllable fast-to-slow light propagation (and vice versa) can be achieved with manipulating the parameter regimes.

## 2. Model and theory

Figure 1 shows the schematic setup that will be studied in this paper, where a hybrid artificial molecule composed of a spherical MNP with radius $R$ and a spherical SQD separated by a distance $d$ in the presence of a strong pump field and a weak probe field. The Hamiltonian of the hybrid SQD-MNP system is given by [7, 23, 39, 41, 46]

$$
H = \hbar \omega_{\text{ex}} \sigma^z + \hbar \omega_{\text{SPP}} a^+ a + \hbar g (\sigma^+ a + a^+ \sigma^-) \\
- \mu E_r (\sigma^+ e^{-i\omega_{\text{ex}}t} + \sigma^- e^{i\omega_{\text{ex}}t}) \\
- \mu E_i (\sigma^+ e^{-i\omega_{\text{iP}}t} + \sigma^- e^{i\omega_{\text{iP}}t}),
$$

(1)

where the first term indicates the Hamiltonian of the SQD with the exciton frequency $\omega_{\text{ex}}$. We consider the SQD is a two-level system (TLS) including the ground state $|g\rangle$ and the single exciton state $|\text{ex}\rangle$ [48, 49] at low temperature, and the two-level exciton can be characterized by the pseudospin $-1$ operators $\sigma^z$ and $\sigma^\pm$.

MNPs can be excited to produce surface-plasmon polaritons (SPPs), which provides an external localized plasmon field [50] and enhances the coherent optical properties of the SQD. There are two representative theoretical descriptions of the exciton-plasmon interaction in the hybrid SQD-MNP system at present, i.e. a semiclassical description [9, 32] and a full quantum description [7, 39–42]. Due to the quantum description for the exciton-plasmon interaction can reveal more novel quantum optical properties that may be applied in quantum processing devices, here we use the quantum treatment to describe the hybrid coupled SQD-MNP system. Quantizing the SPP field in the MNP, then the second term indicates the
Hamiltonian of SPP, where $\omega_{\text{SPP}}$ is the frequency of the SPP mode, and $a$ ($a^\dagger$) is its annihilation (creation) operator. The third term describes the interaction between the SQD exciton and the quantized SPP field with the coupling strength $g$ in the rotating wave approximation [7].

The last two terms in equation (1) indicate the interactions of the probe field and the pump field, respectively. In a rotating frame at the pump field frequency $\omega_z$, we obtain the total Hamiltonian of the system as

$$H = h\Delta_c\sigma^z + h\Delta_{\text{SPP}}a^\dagger a - h\Omega_c (\sigma^+ e^{-i\delta t} + \sigma^- e^{i\delta t}),$$  

(2)

where $\Delta_c = \omega_c - \omega_z$ is the detuning of the exciton frequency and the pump frequency, $\Delta_{\text{SPP}} = \omega_{\text{SPP}} - \omega_z$ is the detuning of the SPP field and the pump frequency, $\Omega_c = \mu E_z / \hbar$ is the Rabi frequency of the pump field, and $\delta = \omega_z - \omega_c$ is the detuning of the probe field and the pump field.

According to the Heisenberg equation of motion and introducing the corresponding damping and noise terms, we derive the quantum Langevin equations [7, 23, 51] as follows

$$\dot{\sigma}^z = -\Gamma_1 (\sigma^z + 1) + ig (\sigma^+ a - \sigma^- a^\dagger)$$

$$+ \Omega_c (\sigma^+ - \sigma^-) + \frac{i\mu E_z}{\hbar} (\sigma^+ e^{-i\delta t} - \sigma^- e^{i\delta t}),$$  

(3)

$$\dot{\sigma}^- = -(i\Delta_c + \Gamma_2) \sigma^- - 2i(ga + \Omega_c) \sigma^- - \frac{2i\mu E_z \sigma^z}{\hbar} e^{-i\delta t} + \hat{\nu},$$

(4)

$$\dot{\hat{a}} = -(i\Delta_{\text{SPP}} + \kappa_{\text{SPP}}/2) \hat{a} + ig \sigma^+ + \hat{\zeta},$$  

(5)

where $\Gamma_1$ ($\Gamma_2$) is the exciton spontaneous emission rate (dephasing rate) and $\kappa_{\text{SPP}}$ is the SPP mode relaxation rate. $\hat{\nu}$ is the $\delta$-correlated Langevin noise operator with zero mean, and $\hat{\zeta}$ is the Langevin force arising from the interaction between the SPP modes and the environment. Using $O = O_0 + \delta O$ ($O = \sigma^z, \sigma^-, a$) equations (3)–(5) can be divided into the steady parts and the fluctuation ones. Substituting the division forms into equations (3)–(5) and setting all the time derivations at the steady parts to be zero, we obtain the steady state solutions of the variables determining the steady-state population inversion ($w_0 = \sigma^0$) of the exciton, which obeys the following equation

$$\Gamma_1 (w_0 + 1) [(\Delta_c^2 + \Gamma_2^2)^2 + \kappa_{\text{SPP}}^2/4]$$

$$+ 2g^2 w_0 (2\Delta_c \Delta_{\text{SPP}} - \Gamma_2 \kappa_{\text{SPP}}) + 4g^4 w_0^2]$$

$$+ 4\Omega_c^2 w_0 \Gamma_2 (\Delta_c^2 + \kappa_{\text{SPP}}^2/4) = 0.$$  

(6)

As all the pump fields are assumed to be sufficiently strong, all the operators can be identified with their expectation values under the mean-field approximation $\langle Q_c \rangle = \langle Q \rangle_e$ [52]. After being linearized by neglecting nonlinear terms in the fluctuations, the Langevin equations for the expectation values

$$\langle \delta \sigma^z \rangle = -\Gamma_1 \langle \delta \sigma^z \rangle - ig [\langle a \rangle \langle \delta a \rangle + \alpha_0 \langle \delta \sigma^+ \rangle - \alpha_0 \langle \delta \sigma^- \rangle - \sigma_0 \langle \delta a^\dagger \rangle]$$

$$+ \Omega_c \langle \delta \sigma^- \rangle + \frac{i\mu E_z}{\hbar} \sigma_0 e^{-i\delta t} - \sigma_0 e^{i\delta t}],$$

(7)

$$\langle \delta \sigma^- \rangle = - (i\Delta_c + \Gamma_2) \langle \delta \sigma^- \rangle - 2ig [w_0 \langle \delta a \rangle$$

$$+ a_0 \langle \delta \sigma^+ \rangle] - 2i\Omega_c \langle \delta \sigma^z \rangle = \frac{2i\mu w_0 E_z}{\hbar} e^{-i\delta t},$$

(8)

$$\langle \delta \hat{a} \rangle = -(i\Delta_{\text{SPP}} + \kappa_{\text{SPP}}/2) \langle \delta \hat{a} \rangle + ig \langle \delta \sigma^+ \rangle,$$  

(9)

which is a set of nonlinear equations and the steady-state response in the frequency domain is composed of many frequency components. To solve these equations, we make the ansatz [53] as $\langle \delta O \rangle = O_e e^{-i\delta t} + O_o e^{i\delta t}$, and substitute them into equations (7)–(9) with ignoring the second-order terms and working to the lowest order in $E_z$ but to all orders in $E_c$, we obtain the linear optical susceptibility as

$$\chi_{\text{eff}}^{(1)}(\omega) = \mu \sigma_+(\omega)/E_c = (\mu^2/h\Gamma_2) \chi^{(1)}(\omega),$$

(10)

and $\chi^{(1)}(\omega)$ is given by

$$\chi^{(1)}(\omega) = \frac{[\Lambda_2 \Pi_3 (\Lambda_1^2 + \Lambda_1^2 \Pi_1) - 2iw_0 \Lambda_1^2] \Gamma_2}{\Lambda_1 \Lambda_3^2 - \Pi_2 \Pi_4 \Lambda_1 \Lambda_2}. $$
where $\beta_1 = ig/(\beta s + \kappa_{spp}/2 - i\delta)$, $\beta_2 = ig/(\beta s + \kappa_{spp}/2 + i\delta)$, $\Pi_1 = \kappa_{spp}/2 - i\delta$, $\Pi_2 = \kappa_{spp}/2 + i\delta$, $\Pi_3 = \kappa_{spp}(\Omega_{spp} - i\delta)$, $\Pi_4 = \kappa_{spp}(\Omega_{spp} + i\delta)$.

Based on the hybrid SQD-MNP system, we can determine the light group velocity as \[ v_g = c/[n + \omega_s(dn/d\omega_s)] \] where $n \approx 1 + 2\pi\chi^{(1)}$. Therefore

\[
c/v_g = 1 + 2\pi\Re\chi^{(1)}(\omega_{s\omega_s}) + 2\pi\omega_s\Re(d\chi^{(1)}/d\omega_s)_{\omega_s=\omega_s},
\] (11)

Obviously, when $Re\chi^{(1)}(\omega_{s\omega_s}) = 0$, the dispersion is steeply positive or negative, and the group velocity is significantly reduced or increased. Thus we define the group velocity index $n_g$ as

\[
n_g = c/v_g - 1 = \frac{c - v_g}{v_g} = \frac{2\pi\omega_s\mu^2}{\hbar \Pi} Re\left(\frac{d\chi^{(1)}}{d\omega_s}\right)_{\omega_s=\omega_s},
\] (12)

where $\Pi = 2\pi\omega_s\mu^2/\hbar$. One can observe the slow light when $n_g > 0$, and the superluminal light when $n_g < 0$ [53].

For illustration of the numerical results, we choose the realistic hybrid InAs SQD-Au MNP complex embedded in dielectric medium with constant permittivity $\varepsilon_0 = 1.8$ and $\varepsilon_r = 2$ [7, 41]. For InAs SQD, we use the realistic experimental parameters [56]: $\kappa_{SQD} \approx 1.0 \text{ GHz}$, $\Gamma_1 = 2\kappa_{SQD}$, $\Gamma_2 = \kappa_{SQD}$, $\mu = e\tau_0$ with $\tau_0 = 1 \text{ nm}$, and $\eta = 9.5$, $\hbar\omega_{spp} = 9 \text{ eV}$, $\hbar\omega_{spp} = 0.07 \text{ eV}$, $\Lambda = 0.25$, and $\nu_f = 1.4 \text{ nm/fs}$. Owing to dissipative losses in the metal, the SPP mode displays a very fast relaxation time [22]. The SPP mode relaxation rate used in the Au MNP is $\kappa_{spp} \approx 3.0 \text{ THz}$, and for the coupling strength between the exciton and SPP field, which depend on the distance $d$ between the SQD and the MNP [7, 41]. Here, the order of magnitudes of the coupling strength $g$ is terahertz (10^{12}Hz). The Rabi frequency $\Omega_c$ of the pump field is $\Omega_c = 10 \text{ GHz}$.

3. Numerical results and discussion

3.1 Case A: the exciton-pump field detuning $\Delta_c = 0$ and the SPP-pump field detuning $\Delta_{spp} = 0$

We first consider the condition of $\Delta_c = 0$ and $\Delta_{spp} = 0$ (i.e. resonance), where the interaction between the SQD and MNP is similar to the J-C Hamiltonian of the standard model. Then, the probe absorption spectra will present symmetric splitting, i.e. vacuum Rabi splitting. Figure 2 displays the imaginary part of the dimensionless susceptibility $\Im(\chi^{(1)})$, which indicates the probe absorption spectra of the probe field as a function of the probe detuning $\Delta_s = \omega_s - \omega_{spp}$ with a series of coupling strengths $g$ under the condition of resonance.

Furthermore, the absorption dip approaches zero around $\Delta_c = 0$ with increasing the coupling strength $g$, which means the input probe field is transmitted to the coupled system without experiencing any absorption. The phenomenon is analogous electromagnetically induced transparency (EIT) [58] in $\Lambda$-type atoms systems, which may indicate the slow light effect based on the hybrid SQD-MNP system. We can further interpret the phenomenon with the destructive quantum interference effect between SPP modes and the beat frequency $\delta$ of the two input laser fields radiating on the SQD. Once the beat frequency $\delta$ approximates to the resonance frequency $\omega_{spp}$ of the SPP field, the SPP modes begin oscillating coherently, which leads to Stokes-like $(\Delta_s = \omega_s + \omega_{spp})$ and anti-Stokes-like $(\Delta_s = \omega_s + \omega_{spp})$ scattering of light from the SQD. When the condition is highly off-resonant, the Stokes-like scattering is strongly suppressed. However, the anti-Stokes field can interfere with the near-resonant probe field and thus modify the probe absorption spectra. As a result, the probe absorption spectra presents zero absorption.

In figure 3(a), we give the real part of the dimensionless susceptibility $\Re(\chi^{(1)})$, i.e. the dispersion of the probe light as a function of $\Delta_c$ with a series of coupling strengths $g$ under $\Delta_c = \Delta_{spp} = 0$. Obviously, with increasing the coupling strength $g$, the dispersion of the probe light varies from negative to positive around $\Delta_c = 0$, which results in positive
group delay or slow light propagation through the system. Therefore, the change of dispersion from a negative to positive slope with manipulating the coupling strength corresponds to the control of slow light propagation. Figure 3(b) plots the group velocity index of the probe laser as a function of the coupling strengths under the cases of $\Delta_c = \Delta_{spp} = 0$. One can see that the group velocity index is positive with the variation of the coupling strengths, which represents the slow light effect.

3.2. Case B: the SPP-pump field detuning $\Delta_{spp} = 0$ under several $\Delta_c \neq 0$

Then, we switch the exciton-pump field detuning $\Delta_c$ from $\Delta_c = 0$ to $\Delta_c \neq 0$, and when the pump field is detuned from the exciton transition (i.e. $\Delta_p \neq 0$), the scenario of the probe absorption becomes completely different. Figure 4 shows the probe absorption spectra as a function of $\Delta_c$ with fixed pump intensity under the condition of $\Delta_{spp} = 0$, which experiences the switch from unsymmetric splitting (i.e. Fano resonance) to symmetric splitting (EIT) to Fano resonance with the change of the detuning $\Delta_c$ from $\Delta_c = -3.0$ THz to $\Delta_c = 3.0$ THz. Unlike EIT presenting a symmetric transparency window, the Fano line profile shows an asymmetry shape caused by the scattering of light amplitude when the condition of observing EIT is not met and an extra frequency detuning is introduced. It is obvious that the exciton-SPP field interaction in the hybrid SQD-MNP system induced resonances has a Fano-like shape that varies with the detuning $\Delta_c$. Particularly, when the pump-resonance detuning is $\Delta_c = 0$, the Fano-like resonance changes into a symmetric Lorentzian-shaped absorption peak (i.e. EIT).

In addition, the evolution of the two sharp unsymmetric Lorentzian peaks change significantly under different detuning $\Delta_c$. It is obvious that the amplitude intensity of the left Lorentzian peaks are reduced while the right peaks are enhanced with changing the detuning $\Delta_c$ from $\Delta_c = -3.0$ THz to $\Delta_c = 3.0$ THz. Different from the condition of the exciton-pump field detuning $\Delta_c = 0$ in figure 2, the probe absorption splits into a doublet where each peak has equal strength presenting symmetrical splitting. However, when $\Delta_c \neq 0$, the absorption peaks corresponding to the splitting are asymmetric, and a prominent avoided crossing phenomenon occurs in the hybrid system [59]. This behavior may be ascribed to the off-resonant coupling between the SQD and the SPP mode. The vacuum Rabi oscillation is direct evidence of the coherent energy exchange between the emitter and the cavity photon field in QED. In this context, the MNP in the hybrid SQD-MNP system plays the role of the cavity in C-QED [47]. The probe absorption splits into two resonances, known as the Autler–Townes (AT) splitting, which is also observed in strongly driven QD systems [60]. Obviously, the absorption spectra can be modified effectively via the off-resonant coupling between the SQD and SPP mode.

Figure 3. (a) The dispersion of the probe light as a function of $\Delta_s$ with a series of coupling strengths $g$ under $\Delta_c = \Delta_{spp} = 0$. (b) The group velocity index $n_g$ of probe laser as a function of the coupling strengths $g$.

Figure 4. The probe absorption spectra as a function of $\Delta_c$ for several detuning $\Delta_c$ with fixed pump intensity under the condition of $\Delta_{spp} = 0$. 
The steep dispersion of the Fano resonance profile promises applications in slow-light devices. Figure 5(a) plots the dispersion of the probe light for several different coupling strengths \( g \), and the dispersion changes significantly. We see that the steep slope around \( \Delta_s = 0 \) is different from in figure 3(a), and in figure 5(b), we investigate the group velocity index \( n_g \) as a function of \( g \) under three different detuning \( \Delta_c \), which presents the slow-fast-slow light in the hybrid system.

3.3. Case C: the exciton-pump field detuning \( \Delta_c = 0 \) under several \( \Delta_{spp} \neq 0 \)

In contrast, we further consider another condition, i.e. \( \Delta_c = 0 \) for several different \( \Delta_{spp} \), to make comparison with Case B. In figure 6, we present a series of probe absorption spectrum as a function of \( \Delta_s \) under the condition of \( \Delta_c = 0 \), which also experiences the switch from Fano resonance to EIT to Fano resonance with the change of the detuning \( \Delta_{spp} \) from \( \Delta_c = -3.0 \) THz to \( \Delta_c = 3.0 \) THz. Similarly, the absorption peaks corresponding to the splitting are also asymmetric two Lorentzian peaks, and a prominent avoided crossing phenomenon also occurs in the hybrid system. However, compared with the situation in figure 4, we find that the amplitude intensity of the left Lorentzian peaks is enhanced while the right is reduced when changing the detuning \( \Delta_{spp} \) from \( \Delta_{spp} = -3.0 \) THz to \( \Delta_{spp} = 3.0 \) THz.

Figure 7(a) plots the dispersion of the probe light with several coupling strengths \( g \), which shows the steep positive slope around \( \Delta_c = 0 \). Figure 7(b) plots the group velocity index \( n_g \) of the probe laser under three different \( \Delta_{spp} \). One can see from that the group velocity index \( n_g \) is positive when the coupling strengths \( g \) varies, which represents the slow light effect. This is a little different from in Case B.

3.4. Case D: the exciton-pump field detuning off-resonant \( \Delta_c \neq 0 \) and the SPP-pump field detuning \( \Delta_{spp} \neq 0 \)

Returning to figures 3, 5 and 7 corresponding to Case A, Case B and Case C, we find that almost only slow light effect can appear in the hybrid system. Although the slow to fast light effect can reach in Case B, the signature is not prominent. So we try to control the two detunings of \( \Delta_c \) and \( \Delta_{spp} \), and we find that the Fano resonances are affected by the two detunings significantly as shown in figure 8. Figures 8(a) and (b) show the variation of the probe absorption spectra under two conditions of \( \Delta_c = 2.0 \) THz and \( \Delta_c = -2.0 \) THz with change the detuning \( \Delta_{spp} \) from \( \Delta_{spp} = -3.0 \) THz to \( \Delta_{spp} = 3.0 \) THz, respectively. Compare figure 8(a) with figure 8(b), we find the amplitude intensity of the left Lorentzian peaks are enhanced while the right peaks are reduced with increasing the detuning \( \Delta_{spp} \) from \( \Delta_{spp} = -3.0 \) THz to \( \Delta_{spp} = 3.0 \) THz.
\[ \Delta c = 2.0 \text{ THz and } \Delta c = -2.0 \text{ THz, respectively. In addition, the probe absorption spectra present mirror symmetry at the two conditions. In figures 8(c) and (d), we display a series of probe absorption spectra at } \Delta_{spp} = 2.0 \text{ THz with changing } \Delta_c \text{ from } \Delta_c = -3.0 \text{ THz to } \Delta_c = 3.0 \text{ THz, respectively. Obviously, the amplitude intensity of the left Lorentzian peaks are reduced while the right peaks are enhanced at these conditions. However, there is a common feature in the four figures of figure 8, i.e. Fano resonance vary significantly and the width of the peaks splitting narrow down.}

We further plot the dispersion of the probe light under different parameter regimes corresponding to figure 8 as shown in figures 9(a)–(d), and the rapid dispersion can result in coherent optical propagation. We then plot the group velocity index \( n_g \) of the probe laser under different detuning regimes. In figure 9(e), we investigate the group velocity index \( n_g \) as a function of } \Delta_{spp} \text{.}
a function of \( g \) for three detuning \( \Delta_{\text{app}} \) under \( \Delta_c = 2.0 \) THz, which presents the conversion from slow to fast to slow light. It is obvious that the conversion of fast to slow light effect is more remarkable at bigger detuning \( \Delta_{\text{app}} \) for fixed detuning \( \Delta_c \). Figure 9(f) shows the group velocity index \( n_g \) of the probe laser versus the coupling strength \( g \) for three different detuning \( \Delta_{\text{app}} \) under \( \Delta_c = -2.0 \) THz. It is obvious that the group velocity index \( n_g \) undergoes an advance to delay corresponding to fast-to-slow light, and the phenomenon is more prominent at small detuning \( \Delta_{\text{app}} \) for fixed \( \Delta_c = -2.0 \) THz. Figure 9(g) plots the group velocity index \( n_g \) versus the coupling strength \( g \) and we give three curves at \( \Delta_c = 1.0 \) THz (the red curve), \( \Delta_c = 2.0 \) THz (the green curve) and \( \Delta_c = 3.0 \) THz (the blue one) under \( \Delta_{\text{app}} = 2.0 \) THz, which indicates the conversion from slow to fast light and fast to slow light, respectively. In figure 9(h) we show the group velocity index \( n_g \) as a function of \( g \) for three detuning \( \Delta_c \) under \( \Delta_{\text{app}} = 2.0 \) THz, which can also obtain the conversion from fast to slow light and vice versa. It is different from in figure 9(e) with figure 9(f), the fast-to-slow light or vice versa can be achieved straightforwardly in the hybrid system.

### 4. Conclusions

We have investigated the Fano resonance and coherent optical propagation properties in the hybrid SQD-MNP system, which includes a SQD driven by two-tone fields coupled to MNP. We found vacuum Rabi splitting, the Fano resonances and their related propagation properties such as fast and slow light effects can be achieved under different parameter regimes, such as the exciton-pump field detuning \( \Delta_c \), the SPP-pump field detuning \( \Delta_{\text{app}} \), and the coupling strengths \( g \). With controlling the detuning of \( \Delta_c \) and \( \Delta_{\text{app}} \), a series of asymmetric Fano line shapes can appear in the probe absorption spectrum. In addition, the fast-to-slow light or vice versa can be achieved in the hybrid system by controlling different detuning regimes. The scheme proposed here may provide potential applications in quantum information processing.

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