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\usepackage{graphicx}

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\usepackage[dvipdfm,colorlinks=true, citecolor=blue, urlcolor=blue ]{hyperref}

\begin{document}

\title{Fluorescence from a quantum dot modified by mesoscopic characteristics in nanoplasmonics}

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\author{\*\*\*\*\*\*\*\*\*\*\*\*\*\*}

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\begin{abstract}

The mesoscopic characteristics of a semiconductor quantum dot (QD) provide new opportunities for the manipulation of light-matter interaction. Here, we theoretically study the QD fluorescence, such as the spontaneous emission and the resonance fluorescence, contained in the mesoscopic QD effects determined by the plasmon-matter interaction in a hybrid QD-metal nanosystem. It is revealed that the QD fluorescence experiences different rotation symmetry over the QD spatial orientation and shows different features compared with that under the dipole approximation when the QD is positioned within the penetration depth of the plasmons into the dielectric. By making use of the QD mesoscopic characteristics, the fluorescence can be controlled, which may have useful application in the development of nanophotonic devices.

\end{abstract}

\pacs{42.50.Pq, 03.67.Bg, 42.50.Dv}

\maketitle

\section{introduction}

Quantum optics has advanced to the stage of experimental measurement and manipulation of individual quantum systems \cite{Haroche2013a,Wineland2013,Wang2015}, where the light-matter interaction plays an important role. With the further development of technology, there is now great interest in exploring new physical mechanisms that enable efficient control of the interaction between individual quantum systems and photon fields. Experimentally, with the reduction of the effective mode volume for the photons, strong and even ultra-strong light-matter coupling have been realized \cite{Wallraff2004,Guebrou2012,Niemczyk2010,Scalari2012}. Recently, a new scheme exploiting the mesoscopic characteristics of QDs has been proposed \cite{Andersen2011}, by means of which the plasmon-matter interaction can be strongly modified.

In most textbooks, the light-matter interaction is described under the dipole approximation (DA) \cite{Scully1997}, which works well in atomic systems where the variation of the field is negligible over the spatial extent of the emitter \cite{Cronin2009,Haroche2013b}. However, once the spatial variation of the field becomes pronounced, such as when the surface plasmon polaritons, or just plasmons, are bounded to the metal surface \cite{Raether1988,Tame2013}, and the emitter is spatially extended, such as a QD several tens of nanometers in size \cite{Lodahl2015}, the validity of the DA is not clear a priori. Experimentally, a large derivation from the dipole theory was observed for QDs in close proximity to a silver mirror \cite{Andersen2011}. Theoretical work indicates that it is the mesoscopic moment, mainly the quadrupolar term, of the QD that invalidates the DA, which has no equivalence in atomic systems \cite{Stobbe2012}. Recently, a microscopic theory of these mesoscopic quantum effects has been provided, which indicates that the mesoscopic moment originates from a circular quantum current density flowing along a curved path inside the QD \cite{Tighineanu2015}. In practice, as they can be tuned by controlling the size and shape of the QD \cite{Johansen2008}, the QD mesoscopic characteristics can be exploited for the development of nanophotonic devices. Up to now, a significantly improved efficiency of the nanoplasmonic single-photon source with a mesoscopic emitter has been theoretically predicted \cite{Andersen2011}. Furthermore, as we know, the radiation behavior of an emitter is determined by its interaction with the field, so that a study of the QD fluorescence modified and controlled by the mesoscopic QD characteristics is necessary and significant.

In this work, based on a model containing the mesoscopic QD effects, we study the spontaneous emission and resonance fluorescence of QDs at different orientations in a hybrid QD-metal nanosystem. Significant deviations from the dipole theory are found when the QD is positioned within the penetration depth of the plasmons into the dielectric. Firstly, due to the cross-interaction between the dipole moment and the mesoscopic moment, the rotation symmetry of the fluorescence over the QD orientation under the DA is destroyed. Secondly, as a result of the cooperative action between the dipole moment and the mesoscopic moment, features of the fluorescence, such as the decay rate of the spontaneous emission and the widths and intensity of the resonance fluorescence spectrum, deviate significantly from that under the dipole theory. By exploiting the mesoscopic QD effects, the fluorescence can be well controlled. The spontaneous decay rate can be either strengthened or weakened, and the resonance fluorescence spectrum can be switched between a single coherent peak and the Mollow-type triplet by varying the mesoscopic moment. Our parameters are all within the values attained already in actual experiments.

Our paper is organized as follows. In Sec. \ref{model}, we show the model and the QD-field interaction beyond the DA. In Sec. \ref{fluorescence}, we study the fluorescence of the QD modified by the mesoscopic QD effects. In Sec. \ref{con}, a summary is given.

\section{Model and QD-field interaction beyond the dipole approximation}

\label{model}

\subsection{Model and Green tensor}

\begin{figure}[tbp]

\includegraphics[width=0.9\columnwidth]{Fig1.eps}

\caption{(Color online) Diagram of a QD with characteristic frequency $\omega\_0$ placed at a distance $z$ above a metal. The dielectric constant of the meda are given by $\varepsilon\_1$ and $\varepsilon\_m$, respectively.} \label{Fig1}

\end{figure}

The system under study is depicted in Fig. \ref{Fig1}: a QD with a resonance frequency of $\omega\_0$ is embedded in gallium arsenide (GaAs) with a refractive index of $n\_1=3.5$ above a dissipative metal. The metal is characterized by a complex dielectric function $\varepsilon\_m (\omega)$ in the Drude model with $\varepsilon\_m (\omega )=\varepsilon \_{\infty }(1-\frac{\omega \_{p}^{2}}{\omega (\omega +i\gamma \_{p})})$, where $\omega \_{p}$ is the bulk plasma frequency, $\varepsilon \_{\infty}$ is the high-frequency limit of the metal dielectric function, and $\gamma \_{p}$ represents the Ohmic loss responsible for the dissipation in the metal. In our work the metal is silver, and the parameters are chosen to be $\omega \_{p}=3.76$ eV, $\varepsilon \_{\infty}=9.6$, $\gamma \_{p}=0.03\omega\_{p}$ in the frequency range of interest \cite{Johnson1972,Tudela2010}.

Conventionally, the electromagnetic field in dispersive and absorbing dielectrics and its coupling to quantum emitters are described by the Green's tensor $\mathbf{G}(\mathbf{r},\mathbf{r}^{\prime };\omega )$ \cite{Dung1998,Dzsotjan2010}, which is rendered as the field evaluated at position $\mathbf{r}$ due to a single point source at position $\mathbf{r}^{\prime}$ with frequency $\omega$. For this, a solution of the Maxwell-Helmholtz wave equation $[\nabla\times\nabla\times-\frac{\omega^{2}}{c^{2}}\varepsilon(\omega)]\mathbf{G}(\mathbf{r},\mathbf{r}^{\prime };\omega )=\mathbf{I}\delta (\mathbf{r}-\mathbf{r}^{\prime })$ is needed, where $\varepsilon(\omega)$ is the relative dielectric function of the medium at frequency $\omega$, and $\mathbf{I}$ is the identity matrix. For general geometries, we have to resort to numerical simulation, like ??FDTD, FEM, or other numerical methods, but for symmetric geometries such as spheres, cylinders, or planes, analytical expressions of the Green's tensor can be obtained \cite{Novotny2006,Dzsotjan2010,Zubairy2014}. In our configuration, in the upper half-space of the interface, the Green's tensor is calculated by the sum of the vacuum Green's tensor and the reflected Green's tensor $\mathbf{G}(\mathbf{r},\mathbf{r}^{\prime },\omega )=\mathbf{G}\_{0}(\mathbf{r},\mathbf{r}^{\prime },\omega )+\mathbf{G}\_{\text{R}}(\mathbf{r},\mathbf{r}^{\prime},\omega )$. For more detail, see the appendix \ref{Green}.

\subsection{QD-field interaction beyond the DA}

The QD-field interaction is described by means of the minimum coupling Hamiltonian \cite{Stobbe2012}. Working in the radiation gauge, the interaction Hamiltonian reads $\hat{H}\_{\text{int}}(\mathbf{r},t)=-\frac{q}{m}\mathbf{A}(\mathbf{r},t)\cdot \mathbf{\hat{p}}$, where $\mathbf{\hat{p}}$ is the canonical momentum operator, $q$ and $m$ are the charge and mass of the electron, respectively, and $\mathbf{A}(\mathbf{r},t)$ is the vector potential of the field. In second quantization, $\mathbf{A}(\mathbf{r},t)$ is expanded as $\mathbf{A}(\mathbf{r},t)=\sum\_{l}\sqrt{\frac{\hbar }{2\omega\_{l}\varepsilon \_{0}}}[\mathbf{A}\_{l}(\mathbf{r})\hat{a}\_{l}e^{-i\omega\_{l}t}+\text{h.c.}]$, where $\mathbf{A}\_l(\mathbf{r})$ is the field spatial distribution function, $\hat{a}\_{l}$ is the field annihilation operator with frequency $\omega\_{l}$, $\varepsilon\_0$ is the vacuum dielectric function and $l=(\mathbf{k},s)$ is the combined wave vector $\mathbf{k}$ and polarization index $s\in(1,2)$. To go beyond the DA, a Taylor expansion of $\mathbf{A}\_{l}\mathbf{(r)}$ to the first order around the center of the QD has to be made \cite{Andersen2011} $\mathbf{A}\_{l}(\mathbf{r})\approx\mathbf{A}\_{l}(\mathbf{r}\_{0})+\mathbf{DA}\_{l}(\mathbf{r})|\_{\mathbf{r=r}\_{0}}\cdot (\mathbf{r-r}\_{0})$, where $\mathbf{DA}\_{l}(\mathbf{r})$ denotes the Jacobian matrix of partial derivatives of $\mathbf{A}\_{l}(\mathbf{r})$. Besides, in the strong confinement regime, the QD can be described as a two-band model with states $| c\rangle$ and $| v\rangle$ representing an electron and a hole in the conduction and heavy valence band, respectively \cite{Stobbe2009}. Thus, in the interaction picture and employing the rotating wave approximation, the interaction Hamiltonian beyond the DA can be expressed as

\begin{equation}

\hat{H}\_{\text{I}}(t)=\hbar \sum\_{l}(g\_{l}e^{i\Delta \_{l}t}\hat{\sigma}\_{-}\hat{a}\_{l}^{\dag }+\text{h.c.}) \label{H-inter0}

\end{equation}

where $\hat{\sigma}\_{-}=|v\rangle \langle c |$, and $\Delta\_l=\omega\_{l}-\omega\_{0}$ is the frequency detuning. The coupling coefficient is governed by

\begin{equation}

g\_{l}=-\frac{q}{m}\sum\_{j,k}(\frac{1}{2\hbar \epsilon \_{0}\omega \_{l}})^{1/2}[(\mu \_{j}+\Lambda \_{j,k}\nabla \_{k})A\_{lj}^{\ast }(\mathbf{r})]\_{\mathbf{r=r}\_{0}} \label{couple}

\end{equation}

where $j$ and $k$ index coordinates $x,y,z$ in the Cartesian coordinates, $\mu\_j=\langle v|\hat{p}\_j|c\rangle $ and $\Lambda\_{j,k}=\langle v|\hat{p}\_{j}r\_k|c\rangle $ denote the point-dipole moment and the first-order mesoscopic moment of the QD, respectively, and $\nabla \_{k}$ represents the partial derivation of $A\_{lj}^{\ast }(\mathbf{r})$ over $k$.

As can be seen, both the dipole moment and the mesoscopic moment (mainly the quadrupolar term) of the QD contribute to its coupling with the field mode. In this process, the former couples to the field distribution function while the latter couples to the gradient of the field distribution function. In atomic systems, $\Lambda\_{j,k}/\mu\_j \sim r\_k \ll 1$ and $\nabla\_k A\_{l,j}^{\ast }(\mathbf{r})|\_{\mathbf{r}=\mathbf{r}\_0} \approx 0$, contributions from the mesoscopic moment can be safely abandoned and the DA works. However, in our model, as the QD is positioned near the surface of a dissipative metal, where the field intensity exponential decays perpendicular to the surface and the emitter is spatially extended, contributions from the mesoscopic moment cannot be ignored and the DA breaks down.

Generally, the light-matter interaction is characterized by the spectral density $J(\omega )=\sum\_{k}g\_{k}^{2}\delta (\omega -\omega \_{k})$ \cite{Weiss2007}. From Eq. (\ref{couple}), including contributions from the mesoscopic moment, it takes the form

\begin{eqnarray}

J (\omega) &=&\frac{q^{2}}{\pi c^{2}\hbar ^{2}\varepsilon\_{0}m^{2}}\sum\_{j,n,j^{\prime },n^{\prime }}\{(\mu \_{j}^{\ast }+\Lambda

\_{j,n}^{\ast }\nabla \_{n}) \label{spectral}\\

&&\times (\mu \_{j^{\prime }}+\Lambda \_{j^{\prime },n^{\prime }}\nabla\_{n^{\prime }}^{\prime })\text{Im}[G\_{j,j^{\prime }}(\mathbf{r},\mathbf{r}^{\prime };\omega )]\}\_{\mathbf{r=r}^{\prime }=\mathbf{r}\_{0}} \nonumber

\end{eqnarray}

where the subscripts $j,n$ ($j^\prime,n^\prime$) index coordinates $x,y,z$ ($x^\prime,y^\prime,z^\prime$), and $G\_{j,j^{\prime }}(\mathbf{r},\mathbf{r}^{\prime };\omega )$ is the $(j,j^\prime)$ element of the Green's tensor $\mathbf{G}(\mathbf{r},\mathbf{r}^{\prime };\omega )$. In the derivation, the relation $\textrm{Im}[\mathbf{G}(\mathbf{r,r}^{\prime };\omega )]=\frac{\pi c^{2}}{2\omega }\sum\_{l }\mathbf{A}\_{l }^{\ast }(\mathbf{r})\otimes\mathbf{A}\_{l }(\mathbf{r}^{\prime })\delta (\omega-\omega \_{l })$ has been utilized \cite{Stobbe2012}, where $\otimes$ denotes the outer product. In the following, we shall study the spectral density for different orientations of the QD.

\subsection{Spectral density for different QD orientations}

We theoretically study the spectral density defined by Eq. (\ref{spectral}) by rotating the QD along an arbitrary axis through an angle $\phi$ to simulate the QD at different spatial orientations. In parallel, the wave functions of the electron, hole and moments of the QD are rotated in the same way. Taking the experimental parameters of ?? et al\cite{Andersen2011}, initially, the moments are $\vec{\mu}=\bar{\mu}\left(\begin{array}{c}1 \\i \\0\end{array}\right)$ and $\mathbf{\Lambda }=\bar{\Lambda}\left(\begin{array}{ccc}0 & 0 & 0 \\0 & 0 & 0 \\1 & i & 0\end{array}\right)$, from which it can be verified that the spectral density is rotation-invariant along the $z$ axis and rotation-equivalent along the other two axes. In our study, we rotate the QD around the $x$ axis as an example.

Defining the rotator $\hat{U}\_{x}(\phi)=e^{-\frac{i}{\hbar }\hat{L}\_{x}\phi}$ with $\hat{L}\_{x}$ being the angular momentum in the $x$ direction, the rotated states of the QD becomes $|\alpha\rangle \_{x}=\hat{U}\_{x}(\phi)|\alpha\rangle $ ($\alpha=c,v$), according to which we obtain the rotated moments $\vec{\tilde{\mu}}(\phi)=\langle c|\hat{U}\_{x}^{\dag }(\phi)\mathbf{\hat{p}}\hat{U}\_{x}(\phi )|v\rangle $ and $\tilde{\mathbf{\Lambda}}(\phi)=\langle c|\hat{U}\_{x}^{\dag }(\phi)\mathbf{\hat{p}r}\hat{U}\_{x}(\phi )|v\rangle $. They take the form

\begin{eqnarray}

\vec{\tilde{\mu}}(\phi) &=&\bar{\mu}\left(

\begin{array}{c}

1 \\

i\cos \phi \\

i\sin \phi

\end{array}%

\right) , \\

\tilde{\mathbf{\Lambda}}(\phi) &=&\bar{\Lambda}\left(

\begin{array}{ccc}

0 & 0 & 0 \\

-\sin \phi & -i\sin \phi \cos \phi & -i\sin \phi \sin \phi \\

\cos \phi & i\cos \phi \cos \phi & i\cos \phi \sin \phi

\end{array}%

\right) .

\end{eqnarray}

Inserting Eq ?? into Eq. (\ref{spectral}), we obtain the spectral density

\begin{equation}

J(\omega,\phi )=J \_{0}(\omega,\phi )+J \_{\text{R}}(\omega,\phi )

\end{equation}

where the first term originates from the free-space field and the second term from the reflected field. Expressed in cylindrical coordinates, they are governed by

\begin{eqnarray}

J\_{0}(\omega ,\phi ) &=&\frac{\omega }{\Phi }\int\_{0}^{\infty }ds\textrm{Re}[A\_{1}(s,\phi )\bar{\mu}^{2}+B\_{1}(s,\phi )\bar{\Lambda}^{\prime 2}\text{]}, \\

J\_{\text{R}}(\omega ,\phi ) &=&\frac{\omega }{\Phi }\int\_{0}^{\infty }ds\textrm{Re}[(A\_{2}(s,\phi )\bar{\mu}^{2}+B\_{2}(s,\phi )\bar{\Lambda}^{\prime 2} \label{ref}\\

&&+B\_{3}(s,\phi )\bar{\mu}\bar{\Lambda}^{\prime })\exp (2is\_{z}k\_{1}z)]. \nonumber

\end{eqnarray}

where $\Phi =(\frac{q^{2}n\_{1}}{8\pi^2 \varepsilon \_{0}m^{2}\hbar^{2}c^{3}})^{-1}$, $s\_{z}=\sqrt{1-s^{2}}$, and $\bar{\Lambda}^{\prime}=k\_{1}\bar{\Lambda}$ with the wavevector $k\_1=n\_1\omega/c$. Perpendicular to the surface, the spatial distribution of the spectral density is determined by $J\_{\text{R}}(\omega,\phi )$, whose integration range $[0,\infty]$ can be divided into two intervals $[0,1]$ and $[1,\infty]$. The former is associated with the reflected radiation modes, which are oscillatory plane waves in space, whereas the latter is associated with the evanescent modes, which exponentially decay away from the metal surface and are bound modes. Furthermore, as we can see, it is a cooperative action between the dipole moment and the mesoscopic moment of the QD that determines the QD-field interaction. The strength coefficients are given as

%{\small{

\begin{eqnarray}

&&A\_{1}(s,\phi ) =\frac{s}{s\_{z}}[(2-s^{2})(1+\cos ^{2}\phi )+2s^{2}\sin

^{2}\phi ], \label{aa}\\

&&A\_{2}(s,\phi ) =\frac{s}{s\_{z}}[(r^{\text{s}}-r^{\text{p}}s\_{z}^{2})(1+\cos ^{2}\phi

)+2s^{2}\sin ^{2}\phi r^{\text{p}}], \nonumber

\end{eqnarray}%}}

and

%{\small{

\begin{eqnarray}

&&B\_{1}(s,\phi ) =\frac{s}{s\_{z}}\{s^{2}[1+7s\_{z}^{2}-(\frac{1}{4}+7s\_{z}^{2})\sin ^{2}\phi]\sin ^{2}\phi \label{bb}\\

&&+ s^{4}(1+\cos ^{2}\phi )\cos ^{2}\phi+(2-s^{2})\sin ^{4}\phi\}, \nonumber\\

&&B\_{2}(s,\phi)=\frac{s}{s\_{z}}\{\frac{s^{2}}{4}[4(r^{\text{s}}-3r^{\text{p}}s\_{z}^{2})-(r^{\text{s}}-11r^{\text{p}}s\_{z}^{2})\sin ^{2}\phi ] \nonumber\\

&&\times\sin^{2}\phi-s\_{z}^{2}(r^{\text{s}}-r^{\text{p}}s\_{z}^{2})\sin ^{4}\phi +s^{4}r^{\text{p}}(1+\cos ^{2}\phi)\cos ^{2}\phi ]\}, \nonumber\\

&&B\_{3}(s,\phi ) =2is[2s^{2}r^{\text{p}}+(2r^{\text{p}}s^{2}-r^{\text{s}}+r^{\text{p}}s\_{z}^{2})\sin ^{2}\phi]\cos \phi. \nonumber

\end{eqnarray}%}}

where $r^{\text{s}}=\frac{s\_{z}-\sqrt{n\_{m1}^{2}-s^{2}}}{s\_{z}+\sqrt{n\_{m1}^{2}-s^{2}}}$ and $r^{\text{p}}=\frac{\varepsilon (\omega )s\_{z}-\varepsilon \_{1}\sqrt{n\_{m1}^{2}-s^{2}}}{\varepsilon (\omega )s\_{z}-\varepsilon \_{1}\sqrt{n\_{m1}^{2}-s^{2}}}$ are the Fresnel reflection coefficients for s-polarized and p-polarized light with the relative dielectric function $n\_{m1}=\sqrt{\varepsilon\_m(\omega)/\varepsilon\_1}$.

Under the DA, only the dipole moment contributes to the spectral density. From Eq. (\ref{aa}) we find that the strength coefficients are functions of $\phi$ with a period of $\pi$, which implies a $\pi$ rotation symmetry of the spectral density over the QD orientation. However, from the definition of $B\_{3}(s,\phi)$ we find that, if the mesoscopic moment cannot be abandoned, this symmetry will be destroyed. Further more, as $B\_{3}(s,\phi)$ is the strength coefficient of the contribution to the spectral density from the cross-interaction between the dipole moment and the mesoscopic moment. Thus, it is this cross-interaction that breaks the rotation symmetry of the spectral density under the DA. In the next section, we will study fluorescence of the QD modified by the mesoscopic QD effects at different QD orientations.

\section{Fluorescence modified by the mesoscopic QD effects} \label{fluorescence}

\subsection{The spontaneous emission}

\begin{figure}[tbp]

\includegraphics[width=\columnwidth]{Fig2.eps}

\caption{(Color online) Orientation dependence of $\Gamma$ at different renormalized separation $\bar{z}$ is plotted under (a) and beyond (b) the DA. A cross-section view is given in (c) at $\bar{z}=0.3$ (blue thick and green dot-dashed lines from (a) and (b)). (d) plots $\Gamma$ vias $\bar {\Lambda}^\prime/\bar{\mu}$ when $\phi=0$ (purple thick line) and $\phi=\pi$ (orange dashed line). Parameters are chosen as $\omega\_0=1.2$ eV, $\Phi=3.0\times10^{9}\bar{\mu}^2$, and $\bar{\Lambda}^{\prime }=0.3\bar{\mu}$ \cite{Andersen2011}.} \label{Fig2}

\end{figure}

Spontaneous emission of the QD can be well described within the Weisskopf-Wigner theory \cite{Scully1997}. Assuming that the exciton is initially excited ($c\_{e}(0)=1$),its population is given by $|c\_{e}(t)|^{2}=\exp [-\Gamma t]$, where the decay rate is

\begin{equation}

\Gamma =2\pi J(\omega\_0,\phi ) \label{decay-rate1}

\end{equation}

In Fig. \ref{Fig2}(a) and Fig. \ref{Fig2}(b) we plot $\Gamma$, for different values of the QD-interface renormalized separation $\bar{z}=z\omega \_{p}/c$, as a function of the QD orientation under and beyond the DA, respectively. As can be seen, we find that $\Gamma$ rapidly attenuates at small values of $\bar{z}$ and oscillates as it increases. This originates from the spatial distribution of the spectral density, as discussed in Eq. (\ref{ref}). In addition, significant deviations are seen at small values of $\bar{z}$, where the DA is broken. Actually, experimental results indicate that the deviations can always be observed when $\bar{z}$ is within the penetration depth of the plasmons into the dielectric \cite{Andersen2011}. In our parameters, it is about $\bar{z}\thicksim2$. For more detail, a cross-section is given in Fig. \ref{Fig2}(c) at $\bar{z}=0.3$. We find that, $\Gamma$ experiences a $\pi$ rotation symmetry over the QD orientation under the DA, while it is destroyed and changed to $2\pi$ once the mesoscopic effects are taken into account. This agrees with our former discussion of the spectral density and originates from the cross-interaction between the dipole moment and the mesoscopic moment of the QD. Furthermore, modified by the mesoscopic QD effects, the values of $\Gamma$ deviate largely from that under the DA: it can be either increased or decreased at different orientations of the QD. This originates from the cooperative action between the dipole moment and the mesoscopic moment and they can be either coherently added or subtracted at different QD orientations. In Fig. \ref{Fig2}(d), we plot $\Gamma$ vias $\bar{\Lambda}^\prime/\bar{\mu}$ at two different orientations of the QD as $\phi=0$ and $\pi$. We find that $\Gamma$ is always strengthened at $\phi=\pi$, while at $\phi=0$, $\Gamma$ can be either coherently subtracted or added at different values of $\bar{\Lambda}^\prime/\bar{\mu}$, which displays an effective control of the decay rate by exploiting the mesoscopic QD effects.

%At last, note that the results at very short distances ($\bar{z}<0.2$) is not given, because the nonradiative processes in the metallic surface is dominant in these region and $\Gamma$ is divergent \cite{Tudela2010}.

\subsection{The resonance fluorescence}

Hereafter, we consider the case in which the QD is resonantly driven by a laser, so that the resonance fluorescence of the QD can be studied. In a frame rotating at the laser frequency $\omega\_\text{L}$, the master equation reads

\begin{eqnarray}

\dot{\rho}(t) &=&-i\Omega [\hat{\sigma}\_{x},\rho (t)] \label{master}\\

&&+\frac{\Gamma}{2}(2\hat{\sigma}\_{-}\rho (t)\hat{\sigma}\_{+}-\hat{%

\sigma}\_{+}\hat{\sigma}\_{-}\rho (t)-\rho (t)\hat{\sigma}\_{+}\hat{\sigma}\_{-})

\nonumber \nonumber

\end{eqnarray}

where $\Omega$ is the Rabi frequency characterizing the strength of the pumping.

The incoherent spectrum of the driven two-level system is defined as $S(\omega )=\frac{1}{\pi }\textrm{Re}[\int\_{0}^{\infty }d\tau e^{i\omega \tau

}\langle \Delta \hat{\sigma}\_{+}(t)\Delta \hat{\sigma}\_{-}(t+\tau )\rangle \_{\text{ss}}]$, where $\Delta \hat{\sigma}\_{\pm }(t)=\hat{\sigma}\_{\pm }(t)-\langle \hat{\sigma}\_{\pm }(t)\rangle \_{\text{ss}}$ with subscript "ss" denoting the steady state. In strong pumping regime ($\Omega\gg\Gamma/4$), the spectrum is characterized by the Mollow-type triplet, which have three peaks centered at $\omega=-\Omega,0,\Omega$ with full-width at half-maximum $\frac{3}{4}\Gamma$, $\frac{2}{4}\Gamma$, and $\frac{3}{4}\Gamma$, respectively, and the spectrum intensity is proportional to $1/\Gamma$. While with a weak pumping ($\Omega<\Gamma/4$), the spectrum is characterized by a single coherent peak at $\omega=0$ \cite{Carmichael2000}. As can be seen, features of the spectrum is determined by $\Gamma$ at a given pumping power $\Omega$, which should be modified by the mesoscopic QD effects as well as the decay rate.

\begin{figure}[tbp]

\includegraphics[width=\columnwidth]{Fig3.eps}

\caption{(Color online) Incoherent spectrum of the QD is plotted at different QD orientations under (a) and beyond (b) the DA. The line with the same type and color takes the same values. The parameters are the same as Fig. \ref{Fig2}(c) with $\Omega=5$ ns$^{-1}$.} \label{Fig3}

\end{figure}

\begin{figure}[tbp]

\includegraphics[width=\columnwidth]{Fig4.eps}

\caption{(Color online) Incoherent spectrum of the QD is plotted vias $\bar{\Lambda}^\prime$ at $\phi=0$ (a) and $\phi=\pi$ (b). Parameters are the same as in Fig. \ref{Fig3}.} \label{Fig4}

\end{figure}

Figure \ref{Fig3} plots the spectrum at different QD orientations under and beyond the DA. Under the DA, the spectrum takes the highest intensity and smallest full-width at half-maximum at $\phi=0,\pi,2\pi$ and the spectrum periodic changes with a period of $\phi$ over the QD orientation. While beyond the DA, the highest intensity and smallest half-width of at full-maximum of the spectrum arises at $\phi=0,2\pi$ and the period of the spectrum changes to $2\phi$. This agrees with the behavior of the spontaneous decay rate as shown in Fig. \ref{Fig2}(c). Especially, in Fig. \ref{Fig3}(b), modified by the mesoscopic QD effects, we find that the Mollow-type triplet disappears and switches to a single peak around $\phi=\pi$, which originates from the strengthen of the decay rate (see Fig. \ref{Fig2}(c)) leading to $\Omega<\Gamma/4$. Further more, in Fig. \ref{Fig4} we plot the spectrum controlled with a variation of $\Lambda^\prime/\bar{\mu}$ at different QD orientations as $\phi=0$ and $\pi$. As can be seen, the spectrum can be well controlled by exploiting the QD mesoscopic characteristics. At $\phi=0$, features of the Mollow triplet can be either strengthened or weakened with a change of the $\Lambda^\prime/\bar{\mu}$. While $\phi=\pi$ features of the Mollow triplet spectrum is always weakened. In large values of $\Lambda^\prime/\bar{\mu}$, the spectrum can be controlled switching from the Mollow triplet to a single coherent peak.

\section{Conclusion}\label{con}

We study the fluorescence, containing the spontaneous emission and the resonance fluorescence, of QD mediated by the QD mesoscopic characteristics in a hybrid QD-metal nanosystem. Modified by the mesoscopic QD effects, large derivations from the dipole theory are viewed. By exploiting the mesoscopic QD effects, fluorescence of the QD can be well controlled. Our studies are within the present experimental states of the art and instructive for the utilization of the QD mesoscopic characteristics in the nanophotonic device developments.

\section\*{ACKNOWLEDGEMENTS}

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\appendix

\section{Green's tensor of planar interface}\label{Green}

For the planar interface considered in our model, the Green's tensor can be calculated analytically \cite{Novotny2006}

\begin{equation}

\mathbf{G}(\mathbf{r},\mathbf{r}^{\prime };\omega )=\mathbf{G}\_{0}(\mathbf{r},\mathbf{r}^{\prime };\omega )+\mathbf{G}\_{\text{R}}(\mathbf{r},\mathbf{r}^{\prime};\omega )

\end{equation}

which is expressed as a sum over the free-space and the reflected field contributions. Here, positions of the field and source point are taken as $\mathbf{r}=(x,y,z)$ and $\mathbf{r}^{\prime }=(x^{\prime },y^{\prime },z^{\prime })$. In the Cartesian coordinates, $\mathbf{G}\_{0}(\mathbf{r},\mathbf{r}^{\prime };\omega )$ and $\mathbf{G}\_{\text{R}}(\mathbf{r},\mathbf{r}^{\prime};\omega )$ are given by

\begin{widetext}

\begin{eqnarray}

\mathbf{G}\_{0}(\mathbf{r},\mathbf{r}^{\prime };\omega ) &=&\frac{i}{8\pi ^{2}%

}\int \int\_{-\infty }^{\infty }dk\_{x}dk\_{y}\frac{e^{i[k\_{x}(x-x^{\prime

})+k\_{y}(y-y^\prime )+k\_{z\_1}|z-z^\prime|]}}{k^{2}k\_{z\_1}}\left(

\begin{array}{ccc}

k\_1^{2}-k\_{x}^{2} & -k\_{x}k\_{y} & \mp k\_{x}k\_{z\_1} \\

-k\_{x}k\_{y} & k\_1^{2}-k\_{y}^{2} & \mp k\_{y}k\_{z\_1} \\

\mp k\_{x}k\_{z} & \mp k\_{y}k\_{z} & k\_1^{2}-k\_{z\_1}^{2}%

\end{array}%

\right),\label{free-green0} \\

\mathbf{G}\_{\text{R}}(\mathbf{r},\mathbf{r}^{\prime };\omega ) &=&\frac{i}{8\pi ^{2}%

}\int \int\_{-\infty }^{\infty }dk\_{x}dk\_{y}\frac{e^{i[k\_{x}(x-x^\prime

)+k\_{y}(y-y^\prime )+k\_{z\_1}(z+z^\prime )]}}{k\_{x}^{2}+k\_{y}^{2}}\times \lbrack

\mathbf{M}^{\text{s}}+\mathbf{M}^{\text{p}}],\label{ref-green0}

\end{eqnarray}%

with $k\_1=|\mathbf{k}\_1|$ and $\mathbf{k}\_1=(k\_{x},k\_{y},k\_{z\_1})$ being the wave vector in medium $1$. Some terms in $\mathbf{G}\_{0}(\mathbf{r},\mathbf{r}^{\prime };\omega )$ have two different signs, which originates from the absolute value of $|z-z^\prime|$. When $z>z^\prime$, the upper sign is applied, and when $z<z^\prime$, the lower sign is applied. However, as $z\equiv z^\prime$ in our work, the different choices are actually equivalent. $\mathbf{G}\_{\text{R}}(\mathbf{r},\mathbf{r}^{\prime };\omega ) $ has been split into the s-polarized part ($\mathbf{M}^{\text{s}}$) and the p-polarized part ($\mathbf{M}^{\text{p}}$). They are given by

\begin{equation}

\mathbf{M}^{\text{s}}=\frac{r^{\text{s}}(k\_x,k\_y)}{k\_{z\_1}}\left(

\begin{array}{ccc}

k\_{y}^{2} & -k\_{x}k\_{y} & 0 \\

-k\_{x}k\_{y} & k\_{x}^{2} & 0 \\

0 & 0 & 0%

\end{array}%

\right) ,\mathbf{M}^{\text{p}}=\frac{-r^{\text{p}}(k\_x,k\_y)}{k\_1}\left(

\begin{array}{ccc}

k\_{x}^{2}k\_{z\_1} & k\_{x}k\_{y}k\_{z\_1} & k\_{x}(k\_{x}^{2}+k\_{y}^{2}) \\

k\_{x}k\_{y}k\_{z\_1} & k\_{y}^{2}k\_{z\_1} &

k\_{y}k\_{z\_1}(k\_{x}^{2}+k\_{y}^{2}) \\

-k\_{x}(k\_{x}^{2}+k\_{y}^{2}) & -k\_{y}(k\_{x}^{2}+k\_{y}^{2}) &

-(k\_{x}^{2}+k\_{y}^{2})^{2}/k\_{z\_1}\end{array}\right).

\end{equation}

\end{widetext}

where $r^{\text{s}}(k\_{x},k\_{y})=\frac{k\_{z\_{1}}-k\_{z\_{2}}}{k\_{z\_{1}}+k\_{z\_{2}}}$ and $r^{\text{p}}(k\_{x},k\_{y})=\frac{\varepsilon\_{2}k\_{z\_{1}}-\varepsilon \_{1}k\_{z\_{2}}}{\varepsilon\_{2}k\_{z\_{1}}+\varepsilon \_{1}k\_{z\_{2}}}$ with $\varepsilon\_{j}$ being the dielectric constant of medium $j(j=1,2)$, represent the Fresnel reflection coefficients for s-polarized and p-polarized light injected from medium $1$ to $2$ respectively. Usually, the medium is unmagnetized and we have taken $\mu\_{j}\equiv\mu\_0=1$, where $\mu\_0$ is the vacuum magnetic susceptibility. $k\_{z\_{2}}$ is the $z$ component of wave vector in medium $2$.

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