Tunable ultranarrow linewidth of a cavity induced by interacting dark resonances

Yandong Peng\textsuperscript{a,b}, Luling Jin\textsuperscript{a,b}, Yueping Niu\textsuperscript{a}\textsuperscript{*} and Shangqing Gong\textsuperscript{b}\textsuperscript{*}

\textsuperscript{a}State Key Laboratory of High Field Laser Physics, Shanghai Institute of Optics and Fine Mechanics, Chinese Academy of Sciences, Shanghai 201800, China; \textsuperscript{b}Graduate University of Chinese Academy of Sciences, Beijing 100049, China

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A scheme for obtaining a tunable ultranarrow linewidth of a cavity due to an embedded four-level atomic medium with double-dark resonances is proposed. It is shown that the steep dispersion induced by double-dark resonances in the transparency window leads to the ultranarrow transmission peak. Compared with the case of a single-dark resonance system, the linewidth can be narrowed even by one order under proper conditions. Furthermore, the position of the ultranarrow peak can be engineered by varying the detuning of the control field.

Keywords: electromagnetically induced transparency; double-dark resonance; optical cavity; absorption; dispersion

1. Introduction

Generally, the narrower the linewidth of a cavity the higher the quality of resonator required. An electromagnetically induced transparency (EIT) \cite{1,2} medium placed in an ordinary cavity can significantly narrow the cavity linewidth, known as intracavity EIT \cite{3}. In experiments, cavity linewidth narrowing has been observed \cite{4}. Although the embedded medium spoils the optical quality factor of the resonator, we can obtain a narrower linewidth with a less perfect cavity. Xiao and co-workers \cite{5} first experimentally demonstrated cavity-linewidth narrowing by means of EIT with hot atomic-Rb vapor in an optical ring cavity. After inserting the atomic vapor, the finesse of the cavity reduces from 100 to 51, but the cavity linewidth is narrower by a factor 7 than the empty-cavity linewidth. Later, in a cold atomic-Rb system, Zhu et al. \cite{6} observed a cavity transmission spectrum with a central peak representing the intracavity dark state and two vacuum Rabi sidebands. Very recently, intracavity EIT and polariton resonance have been observed in a Doppler-broadened medium \cite{7}. The above experiments are based on a single-dark EIT system and the narrow transmission peak lies at the resonance frequency of the selected atomic medium. Large normal dispersion and almost vanishing absorption in the transparency window contribute to the narrow central peak of the cavity transmission. Many other works have been carried out on the cavity+atom system, such as optical bistability and multistability \cite{8–12}, cavity linewidth broadening \cite{13} and cavity transmission in strong-coupling regime \cite{14}, quantum information and memory \cite{15–17}, photon-photon interaction \cite{18}, slow light \cite{19,20}, and so on.

It is well known that dark resonances are the basis of EIT and coherent population trapping, etc. Double-dark resonance was first studied by Lukin et al. \cite{21}. Coherent coupling fields dress atomic states and multiple quantum superposition states appear. Interference for transitions between dressed states and the bare state leads to the splitting of dark states, so-called double-dark resonance. Later, a quantum-interference phenomenon induced by interacting dark resonances was observed in experiments \cite{22} and different schemes of double-dark resonances were explored \cite{23–28}. Some of us have carried out work based on double-dark resonances, such as high efficiency four-wave mixing \cite{29}, enhanced Kerr effect \cite{30}, atom localization \cite{31}, and so on. In this paper, we propose an efficient scheme for getting a tunable ultranarrow cavity transmission by interacting dark resonance. In a tripod configuration, an additional transition to the \( \Lambda \)-type EIT system by another control field causes the occurrence of two distinct dark states. Interacting dark states makes one of the transmission peaks much narrower. By proper tuning of the control fields, the ultranarrow spectrum can be one order of magnitude narrower than that of a single-dark system. Moreover, its position can be controlled by adjusting the detuning of the control field. We also notice that, for single resonance, the central peak of the cavity transmission will become broader when the detuning of
the control field is not zero, while, for double resonance, we could get the ultranarrow linewidth even if the control-field detuning exists. Therefore, the intracavity double-dark system has a tunable ultranarrow linewidth.

2. Model and basic equations

The atomic system under consideration is plotted in Figure 1. Here, \( |0\rangle \) is an exited state and \(|1\rangle, |2\rangle, |3\rangle \) are three Zeeman ground sub-levels. A weak probe field and two control fields with Rabi frequencies \( g, \Omega_1 \) and \( \Omega_2 \) couple the transitions \(|1\rangle - |0\rangle\), \(|2\rangle - |0\rangle\) and \(|3\rangle - |0\rangle\), respectively. The same model has already been used to study slow light, linear susceptibility and enhanced nonlinear optics [24], sub-Doppler and sub-normal absorption line [32], adiabatic pulse propagation [32], quantum [33] and optical [34] switching, and others.

The dynamics of the system can be described using the density-matrix approach as

\[
\dot{\rho} = -\frac{i}{\hbar}[H, \rho] + \mathcal{L}\rho,
\]

where \( H \) is the Hamiltonian of the system and \( \mathcal{L}\rho \) denotes relaxation terms. In the interaction picture, the Hamiltonian of the system in an appropriate rotating regime can be written as

\[
H = \hbar \begin{pmatrix}
0 & -g & -\Omega_1 & -\Omega_2 \\
-g & -\Delta_p & 0 & 0 \\
-\Omega_1 & 0 & -\Delta_1 & 0 \\
-\Omega_2 & 0 & 0 & -\Delta_2
\end{pmatrix}.
\]

And, the Liouvillian matrix \( \mathcal{L}\rho \) is given by

\[
\mathcal{L}\rho = \begin{pmatrix}
\Delta_p - \gamma_1 & (\gamma_1 - \gamma_2 - \gamma_3)\rho_{00} & -\gamma_{01}\rho_{01} & -\gamma_{02}\rho_{02} & -\gamma_{03}\rho_{03} \\
-\gamma_{01}\rho_{10} & \gamma_1\rho_{00} - \gamma_{12}\rho_{12} & -\gamma_{13}\rho_{13} & -\gamma_2\rho_{02} & -\gamma_3\rho_{03} \\
-\gamma_{02}\rho_{20} & -\gamma_{21}\rho_{21} & \gamma_2\rho_{02} - \gamma_3\rho_{03} & -\gamma_1\rho_{12} & -\gamma_3\rho_{13} \\
-\gamma_{03}\rho_{30} & -\gamma_{31}\rho_{31} & -\gamma_{32}\rho_{32} & \gamma_3\rho_{03} & -\gamma_1\rho_{13} \\
\end{pmatrix}.
\]

where \( \Delta_p = \omega_p - \omega_0 + \omega_1 \), \( \Delta_1 = \omega_1 - \omega_0 + \omega_2 \), and \( \Delta_2 = \omega_2 - \omega_0 + \omega_3 \) denote the detunings of the corresponding fields. \( \omega_j \) (\( j = 0 \sim 3 \)) is the atomic eigenfrequency and \( \gamma_j \) represents the spontaneous decay rate from state \( |0\rangle \) to state \( |j\rangle \).

We can easily get the density matrix equations of the system as follows [24]:

\[
\dot{\rho}_{00} = -\left( \gamma_1 + \gamma_2 + \gamma_3 \right)\rho_{00} - ig(\rho_{10} - \rho_{01}) - i\Omega_1(\rho_{20} - \rho_{02}) - i\Omega_2(\rho_{30} - \rho_{03}),
\]

\[
\dot{\rho}_{11} = \gamma_1\rho_{00} + \gamma_{21}\rho_{22} + \gamma_{31}\rho_{33} - ig(\rho_{10} - \rho_{01}),
\]

\[
\dot{\rho}_{22} = \gamma_2\rho_{00} - \gamma_{21}\rho_{22} + \gamma_{12}\rho_{12} - i\Omega_1(\rho_{20} - \rho_{02}),
\]

\[
\dot{\rho}_{33} = \gamma_3\rho_{00} - (\gamma_{31} + \gamma_{32})\rho_{33} - i\Omega_2(\rho_{30} - \rho_{03}),
\]

\[
\dot{\rho}_{10} = -\Gamma_{10}\rho_{10} - ig(\rho_{10} + \rho_{11}) + i\Omega_1(\rho_{12} + \rho_{21}) + i\Omega_2(\rho_{20} + \rho_{30}),
\]

\[
\dot{\rho}_{20} = -\Gamma_{20}\rho_{20} - i\Omega_1(\rho_{00} + \rho_{12}) + i\Omega_1(\rho_{12} + \rho_{20}) + i\Omega_2(\rho_{20} + \rho_{30}),
\]

\[
\dot{\rho}_{30} = -\Gamma_{30}\rho_{30} - i\Omega_2(\rho_{02} + \rho_{30}) + i\Omega_1(\rho_{12} + \rho_{20}) + i\Omega_2(\rho_{20} + \rho_{30}),
\]

\[
\dot{\rho}_{12} = -\Gamma_{12}\rho_{12} - ig(\rho_{12} - \rho_{21}) + i\Omega_1(\rho_{12} + \rho_{21}) + i\Omega_2(\rho_{20} + \rho_{30}),
\]

\[
\dot{\rho}_{13} = -\Gamma_{13}\rho_{13} - ig(\rho_{13} + \rho_{31}) + i\Omega_1(\rho_{13} - \rho_{31}) + i\Omega_2(\rho_{20} + \rho_{30}),
\]

\[
\dot{\rho}_{23} = -\Gamma_{23}\rho_{23} - i\Omega_1(\rho_{03} + \rho_{32}) + i\Omega_2(\rho_{20} + \rho_{30}),
\]

with \( \rho_{kj}^* = \rho_{jk} \) and the closure relation \( \sum_j \rho_{jj} = 1 \) (\( j,k \in \{0,1,2,3\} \)). Here \( \gamma_{jk} \) (\( j \neq k \)) are the relaxation rates of the respective coherences between state \(|j\rangle\) to \(|k\rangle\). \( \Gamma_{jj} = \gamma_{jj} - \Omega_{jj}^2 \), \( \Gamma_{kj} = \gamma_{kj} - i\Omega_{kj} \).

Since a weak probe field couples the transition \(|1\rangle - |0\rangle\), its absorption and dispersion properties are determined by the susceptibility of the intracavity atomic medium. Considering the initial condition (i.e. \( \rho_{11} = 1 \)), weak field approximation and ignoring the decay between ground states (\( \gamma_{13} = \gamma_{12} = \gamma_{23} = 0 \)), the steady state linear susceptibility [23] is given by

\[
\chi(\Delta_p) = -4\pi N N_{\mu 0} \left\langle \frac{ab(\Lambda_p, ab - \Omega_1^2b - \Omega_2^2a - i\gamma_0^2 ab)}{(\Lambda_p, ab - \Omega_1^2b - \Omega_2^2a)^2 + \gamma_1^2 a^2 b^2} \right\rangle,
\]

where \( N \) is the medium density, \( a = \Delta_p - \Delta_1 \) and \( b = \Delta_p - \Delta_2 \). It is easy to see that when \( a = 0 \) or \( b = 0 \), the atomic susceptibility is zero, which means a transparency window at \( \Delta_p = \Delta_1 \) or \( \Delta_p = \Delta_2 \). Moreover, the additional control field to a \( \lambda \)-type system leads to new interference effects and splits the transparency window into two if \( \Delta_1 \neq \Delta_2 \) [23]. By close inspection of Equation (5), we also find that the dependence of these transparency windows on the detunings of the control fields provides the feasibility of manipulating the atomic response at different frequencies.

The linear susceptibilities have been studied in [23,24]. Here, some similar results are plotted in Figure 2 and all parameters are scaled by the decay rate \( \gamma_1 \) for simplicity. In the following, we give a brief
dramatic change of the atomic susceptibility. When the coupling field makes the original dark resonance split, interpret our later results. In essence, the additional detuning $D$ can be expressed as [3]

$$S(\omega) = \frac{t^2}{1 + r^2 \kappa^2 - 2r \kappa \cos[\Phi(\omega)]},$$

where $t$ and $r$ are the transmissivity and the reflectivity of both the input and the output mirrors ($t^2 + r^2 = 1$), and for simplicity, we assume that mirror 3 has 100% reflectivity. $\Phi(\nu) = \omega_c(L + l\nu')$ is the total phase shift, and $\kappa = \exp(-\omega l\nu'/c)$ is the medium absorption per round trip. The real part of the atomic susceptibility brings dispersion and additional phase shift, and the imaginary part introduces absorption leading to the attenuation of the field’s amplification. Of course, the solution of the cavity transmission can also be obtained under the small-gain approximation [7,14].

To the single-dark system in resonance, the EIT frequency $\omega_{0s}$ corresponds to the atomic transition frequency $\omega_{01}$. If an empty-cavity resonance frequency $\omega_c$ is sufficiently close to the EIT frequency $\omega_{0s}$, the resonance frequency $\omega_r$ of the cavity + atom system will be pulled to this value according to the relation [3]

$$\omega_r = \frac{1}{1 + \eta} \omega_{0s} + \frac{\eta}{1 + \eta} \omega_{01},$$

where $\eta = \omega_{0s}(L/l)(\partial \nu'/\partial \omega_c)$ denotes a frequency-stabilization coefficient. While the present medium is the double-dark system, there are two EIT frequencies, $\omega_{0d} = \omega_{01} + \Delta_{1,2}$, and then the medium around two transparency windows will pull the resonance frequency $\omega_{rd}$ to their respective EIT frequencies $\omega_{rd}$. The frequency pulling effect depends on $\eta$ which is decided by the dispersion. The larger dispersion, the larger frequency pulling effect, and the better frequency-stabilization effect. For the double-dark system, since the dispersion in the narrow transparency window is larger than that of the single-dark system, the cavity-atom resonance frequency is more strongly pulled to its EIT frequency. We could then get a better frequency-stability effect.

3. Results and discussion

We consider the atomic medium within a vapor cell of length $l$ in an optical ring cavity of length $L$, as shown in Figure 3. The cavity field serves as the probe field and we adjust the cavity field to be resonant with the atomic transition, i.e. the cavity detuning is zero. The control field is injected into the cavity with a polarization beam-splitter and co-propagates through the vapor cell (not circulating in the cavity) with the cavity field, which forms the two-photon Doppler-free configuration [35]. The cavity transmission can be expressed as [3]

$$S(\omega) = \frac{t^2}{1 + r^2 \kappa^2 - 2r \kappa \cos[\Phi(\omega)]},$$

where $t$ and $r$ are the transmissivity and the reflectivity of both the input and the output mirrors ($t^2 + r^2 = 1$), and for simplicity, we assume that mirror 3 has 100% reflectivity. $\Phi(\nu) = \omega_c(L + l\nu')$ is the total phase shift, and $\kappa = \exp(-\omega l\nu'/c)$ is the medium absorption per round trip. The real part of the atomic susceptibility brings dispersion and additional phase shift, and the imaginary part introduces absorption leading to the attenuation of the field’s amplification. Of course, the solution of the cavity transmission can also be obtained under the small-gain approximation [7,14].

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$$\omega_r = \frac{1}{1 + \eta} \omega_{0s} + \frac{\eta}{1 + \eta} \omega_{01},$$

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Here, we focus on the transmission peak in the narrow transparency window. It is instructive to examine the modified cavity linewidth [3,5]. The rate of the cavity linewidth $\Delta \omega_r$ of the double-dark system to that of the single-dark system $\Delta \omega_s$ reads

$$\frac{\Delta \omega_r}{\Delta \omega_s} = \frac{1 + \omega_{0d}(l/2L)(\partial \chi''/\partial \omega_r)}{1 + \omega_{0d}(l/2L)(\partial \chi''/\partial \omega_r)}$$  \hfill (8)

To the atomic-Rb system, $\omega_{0s}$ and $\omega_{0d}$ are the order of $10^{14}$ Hz [4,37], then it is easy to see that the cavity linewidth is in roughly reverse proportion to the dispersion. When the intracavity atomic system is prepared in the double-dark states, the interacting dark states lead to a narrow transparency window. Therefore, the dispersion in the narrow transparent window is much larger than that of the single-dark system [23]. Then the cavity transmission peak could be much narrowed and the ultranarrow linewidth of the cavity appears. Moreover, since the detunings of the control fields determine the positions of the transparency windows, the position of the ultranarrow transmission peak could be controlled just by changing the detuning of the weak control field.

In the following, we present the numerical simulation of the cavity transmission spectrum in Figure 4. First, let us see the ultranarrow spectrum induced by the interacting double resonances. When $\Omega_1 = 2$, $\Omega_2 = 0.3$ and $\Delta_1 = -\Delta_2 = -1$, we get a sharp narrow transmission peak (see Figure 4(a)). In this case the value of expression (8) is 0.083, which means that the cavity linewidth is narrowed by one order. Furthermore, with a different detuning of the weak control field, the position of the ultranarrow transmission peak is changed (see Figure 4(b)). Second, when the intensities of two control fields are equal, the right peak is broadened because here the steep slope of the atomic susceptibility gets gentle (see Figure 2(c)). Then we get a symmetrical double-peak transmission (see Figure 4(c)). Finally, for comparison, we plot the transmission of the cavity with the single-dark atomic system in Figure 4(d). Obviously, its linewidth is much broader than that of the ultranarrow spectrum. Physically, the double-dark states exist in the tripod system and the interacting dark resonances cause the steep dispersion leading to the ultranarrow spectrum.

In addition, due to the system’s symmetry, we could get the tunable ultranarrow transmission peak on the left of the resonance frequency by tuning the intensity and detuning of the other control field. Then we can control the ultranarrow spectrum of the cavity for both the red and blue detunings, which means more manipulability at a broad frequency range.

4. Conclusion

In conclusion, we proposed a feasible scheme for the ultranarrow linewidth of a cavity with the double-dark system. The results show that the interacting dark resonances narrow the transmission spectrum dramatically. By properly tuning the intensity of the control field, the cavity linewidth could be narrowed one order smaller than that of the single-dark system. Moreover, by adjusting the detuning of the control field, the position of the ultranarrow transmission peak could be controlled. Close inspection of the atomic susceptibility obviously shows that the additional control field causes the splitting of the dark states, and the interacting dark resonances lead to the steep dispersion which, finally, results in the fine spectral features.

In experiments, the tripod atomic system has already been used for beam-splitter [36] and sub-Doppler and sub-natural linewidth absorption spectrum [37], which means it is feasible to achieve a tunable ultranarrow cavity linewidth. This scheme may have potential applications in high-resolution spectroscopy, laser frequency stabilization, optical magnetometry and in studying cavity-QED effects.

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