Microwave assisted transparency in an M-system

Nawaz Sarif Mallick, Tarak N Dey and Kanhaiya Pandey

Department of Physics, Indian Institute of Technology Guwahati, Guwahati, Assam 781039, India

E-mail: kanhaiyapandey@iitg.ernet.in

Received 31 May 2017, revised 26 July 2017
Accepted for publication 5 September 2017
Published 22 September 2017

Abstract
In this work we theoretically study a five-level M-system whose two unpopulated ground states are coupled by a microwave (MW) field. The key feature which makes the M-system more efficient in comparison to a routinely studied closed loop Λ-system is the absence of MW field induced population transfer even at high intensities of the latter. The limitation of closed loop Λ systems due to MW induced population redistribution among the ground states, which reduces the atomic coherences, can be overcome in the M-system. We examine lineshape of probe absorption as a function of its detuning in the presence of both control and MW fields. The MW field facilitates the narrowing of the probe absorption lineshape in M-systems which is in contrast to closed loop Λ-systems. Hence this study opens up a new avenue for atom-based phase-dependent MW magnetometry.

Keywords: EIT, M-system, microwave

(The some figures may appear in colour only in the online journal)

1. Introduction

Recently there are efforts towards atom-based microwave (MW) electrometry and magnetometry due to their high reproducibility, accuracy, spatial resolution and stability [1–5]. This is based upon the phenomenon of electromagnetically induced transparency (EIT) [6, 7] in which the absorption property of a probe laser is altered in the presence of control lasers in a multilevel system. The basis of EIT is the induced or transfer of coherence (TOC) between the levels which are not directly driven by optical control fields. The EIT has been explored extensively in the three level Λ [8–14], V [15–17] and Ξ [18–24] systems. It has also been investigated beyond three level systems such as in doubly driven V [25], N [26–30], Y [31], inverted Y [32], ΞΛ [33, 34], tripod [35, 36] and doubled tripod [37] systems. The role of various TOCs for a general N-level atomic system has been well studied [38]. This phenomenon has been paid a great deal of attention due to its potential applications in a wide variety of fields like controlling the group velocity of light [11, 26], coherent storage and retrieval of light [13, 14], high-resolution spectroscopy [18, 23], studying photon–photon interaction via a Rydberg blockade [19, 20], photon transistor for quantum optical information processing [21, 24, 39], etc.

The induced coherence of a Λ system can be drastically modified by applying a low frequency field which couples two ground levels. Two optical fields along with a lower level (LL) coupling field form a closed loop three level Λ system. The relative amplitude, detuning and phase differences between the LL field and two optical fields can play an important role in significantly manipulating the optical property of the system such as dispersion, absorption and nonlinearity [40–47].

In the above mentioned closed loop Λ system, population transfer is unavoidable in the presence of high power MW field as it acts on one of the populated states. The MW induced population transfer between metastable states leads to imperfect transparency which limits many of the EIT based applications [48]. Here we study an M-system wherein two unoccupied ground states are coupled with a MW field and hence there is no population transfer with this MW Rabi frequency. This MW field also enhances the generated atomic coherence at the probe transition manifold which can not be found in generic three level systems. This enhancement of atomic coherence is a result of the interference between excitation pathways from the different atomic states.

The paper is organized as follows: in the next section, we introduce the physical model and basic equations of motion...
for an M-system by using a semiclassical theory. In section 2.2, an approximate analytical expression for linear susceptibility of a weak probe field is derived using the perturbative approach under three-photon resonance condition. After deriving the analytical form of the atomic coherences, we compare it with the full numerical solution. In section 3, we first explain the lineshape of the probe laser absorption as a function of probe detuning for various MW strengths and phases. Finally M-systems are compared to closed loop Λ-systems both having ground states coupled by MW field.

2. Theoretical formulation

2.1. Model configuration

We consider a five-level atomic M-system as shown in figure 1(a). The electric field associated with the lasers driving the transition $|i\rangle \rightarrow |i+1\rangle$ is $E_{i,i+1} e^{i \phi_{i,i+1}^L}$. The $E_{i,i+1}$ is amplitude, $\omega_{i,i+1}$ is the frequency and $\phi_{i,i+1}^L$ is the phase. We define Rabi frequency $\Omega_{i,j} = d_{ij} E_{i,j} \omega_{i,j}^L / \hbar$ for the transition $|i\rangle \rightarrow |j\rangle$ having the dipole moment matrix element $d_{ij}$ driven by a laser with electric field amplitude $E_{i,j}$ and the phase $\phi_{i,j}^L$. The $\Omega_{12}$ is the Rabi frequency of the probe laser driving $|1\rangle \rightarrow |2\rangle$ and shown with solid blue arrow in figure 1. The Rabi frequencies of the control lasers driving the transitions $|2\rangle \rightarrow |3\rangle$, $|3\rangle \rightarrow |4\rangle$ and $|4\rangle \rightarrow |5\rangle$ are denoted by $\Omega_{23}$, $\Omega_{34}$ and $\Omega_{45}$ respectively and shown with solid red arrows. The lower level states $|3\rangle$ and $|5\rangle$ is coupled by a MW field with Rabi frequency $\Omega_{35}^{MW}$ as shown with solid green arrow. The M-system can be realized in $^{87}$Rb at the D1 line using the two hyperfine levels $F_g = 1$ and $F_g = 2$ of the ground state $5S_{1/2}$ and the hyperfine level $F_e = 1$ of excited state $5P_{1/2}$ as shown in figure 1(b). The decay rate of the excited state $5P_{1/2}$ is $2\pi \times 6$ MHz. We apply a large magnetic field to split the magnetic sublevels to keep away the unwanted sublevels from resonance and to define the quantization axis. The sign and value of the Lande $g$ factor for the chosen hyperfine is suitable to split and keep away the unwanted transitions. At 200 G applied magnetic field the splitting of the magnetic sublevels for $F_e = 1$ is $2\pi \times 46$ MHz and splitting for $F_g = 2$ and $F_g = 1$ is $2\pi \times 140$ MHz as shown in figure 1(c). We consider the probe to be $\sigma^+$ polarized and driving the transition $|F_g = 2; m_f = 0\rangle \rightarrow |F_e = 1; m_f = 0\rangle$. The $\pi$ polarized control laser is driving the $|F_g = 2; m_f = 0\rangle \rightarrow |F_e = 1; m_f = 1\rangle$ and $|F_g = 2; m_f = 0\rangle \rightarrow |F_e = 1; m_f = -1\rangle$ transitions as shown in figure 1(c).

In the absence of the control lasers all the magnetic sublevels of the hyperfine states $F_g = 1$ and $F_g = 2$ will be equally populated. However when the control lasers are applied, the population in the ground states connected with the control lasers will deplete. In the steady state there will be no population left in the $|F_g = 1; m_f = 0\rangle$ and $|F_g = 2; m_f = 0\rangle$ states as shown in figure 1(c).

The experimental set-up is shown in figure 1(d). In this all the optical and MW fields are propagating perpendicular to the atomic beam and hence does not involve any Doppler shift and broadening. The probe laser ($\Omega_{12}$) and the two control lasers, ($\Omega_{34}$ and $\Omega_{45}$) are mixed by a polarizing beam splitter (PBS) and are propagating along the z-direction. The polarization of the probe laser is $\sigma^+$ and it is $\sigma^-$ for the two control lasers, $\Omega_{34}$ and $\Omega_{45}$ w.r.t the applied magnetic field of 200 G along the z-direction as shown in figure 1(d). The opposite circular polarization for control and the probe lasers are set by a $\lambda/4$ waveplate just before entering the glass cell. After coming out of the glass cell the probe laser is separated from the two control lasers using a $\lambda/4$ waveplate and a PBS. The $\pi$-polarized control laser ($\Omega_{35}$) and the MW field ($\Omega_{35}^{MW}$) are propagating along y-direction having the polarization along the z-direction. Instead of a Rb atomic beam we can also use the cold atoms or a nano-cell to avoid Doppler broadening [49, 50].

The unperturbed atomic Hamiltonian can be written as

$$H_0 = \sum_{j=1}^{5} \hbar \omega_j |j\rangle \langle j|,$$

where $\hbar \omega_j$ is the energy of the $|j\rangle$ state. Under the action of five coherent fields, the interaction Hamiltonian of the system in the dipole approximation is given by

$$H_I = \frac{\hbar}{2} \left[ e^{i \omega_{01} t} + e^{-i \omega_{01} t} |1\rangle \langle 2| + e^{i \omega_{12} t} + e^{-i \omega_{12} t} |2\rangle \langle 3| + e^{i \omega_{23} t} + e^{-i \omega_{23} t} |3\rangle \langle 4| + e^{i \omega_{34} t} + e^{-i \omega_{34} t} |4\rangle \langle 5| + e^{i \omega_{45} t} + e^{-i \omega_{45} t} |5\rangle \langle 3| \right].$$
Hence the total Hamiltonian will be \( H = H_0 + H_r \). We use suitable unitary transformation to eliminate the explicit time dependent part in the Hamiltonian. In rotating wave approximation, the effective interaction Hamiltonian can be expressed as

\[
\hat{H}_I = \hbar \left[ \frac{1}{2} [J]_1 \langle 1 - \delta_{12}^2] \langle 2 | - (\delta_{12} - \delta_{23}) | 3 \rangle \langle 3 | - (\delta_{12} - \delta_{23} + \delta_{34}) | 4 \rangle \langle 4 | - (\delta_{12} - \delta_{23} + \delta_{34} - \delta_{45}) | 5 \rangle \langle 5 | \right] + \frac{\Omega_{12}}{2} \left[ | 2 \rangle \langle 3 | + | 3 \rangle \langle 2 | \right] + \frac{\Omega_{34}}{2} \left[ | 4 \rangle \langle 5 | + | 5 \rangle \langle 4 | \right] + \frac{\Omega_{45}}{2} \left[ e^{i\Delta_{45}t} + e^{i\Delta_{34}t} \right] | 3 \rangle \langle 5 | + \text{c.c.},
\]

(3)

where \( \delta_{12} = \omega_{12} - (\omega_1 - \omega_2) \), \( \delta_{23} = \omega_{23} - (\omega_2 - \omega_3) \), \( \delta_{34} = \omega_{34} - (\omega_3 - \omega_4) \), \( \delta_{35} = \omega_{35} - (\omega_3 - \omega_5) \) and \( \Delta_{34} = \omega_{34} - (\omega_3 - \omega_4) \) known as detunings of the lasers for respective transitions. Now we impose the condition \( \delta_{34} + \delta_{35} - \delta_{34} = 0 \), so that the time dependence is completely eliminated from the effective interaction Hamiltonian. To explore the dynamics of the atomic system, we use the density matrix equations where radiative relaxation of the atomic states are included. The atomic density operator \( \rho \) obey the following equations

\[
\dot{\rho} = -i \hbar \left[ \hat{H}_I, \rho \right] - \frac{1}{2} [\Gamma, \rho],
\]

(4)

where \( \Gamma \) is the relaxation matrix [51].

2.2. Perturbative analysis

To study the response of the medium, we numerically solve the density matrix equations at steady state condition. We derive an analytical expression for the probe field response that is correct to first order in probe field approximation and exact for the all order in control fields. In this weak approximation, there will be no population transfer and hence the evolution of the population i.e. diagonal terms of the density matrix such as \( \rho_{11}, \rho_{22}, \rho_{33}, \rho_{44} \) and \( \rho_{55} \) can be ignored with \( \rho_{11} \approx 1, \rho_{22} \approx 0, \rho_{33} \approx 0, \rho_{44} \approx 0 \) and \( \rho_{55} \approx 0 \). Time evolution for the coherences i.e. off-diagonal terms of the density matrix will be given by following set of equations

\[
\dot{\rho}_{12} = -\gamma_{12} \rho_{12} + \frac{i}{2} \Omega_{12}^* \rho_{13}^* + \frac{i}{2} \Omega_{12} \rho_{34},
\]

\[
\dot{\rho}_{13} = -\gamma_{13} \rho_{13} - \frac{i}{2} \Omega_{12}^* \rho_{12} + \frac{i}{2} \Omega_{12} \rho_{23} + \frac{i}{2} \Omega_{34} \rho_{13} + \frac{i}{2} \Omega_{34}^* \rho_{14} + \frac{i}{2} \Omega_{13} \rho_{15} + \frac{i}{2} \Omega_{13}^* \rho_{16},
\]

\[
\dot{\rho}_{14} = -\gamma_{14} \rho_{14} + \frac{i}{2} \Omega_{13} \rho_{15} + \frac{i}{2} \Omega_{13}^* \rho_{16} + \frac{i}{2} \Omega_{14} \rho_{24} + \frac{i}{2} \Omega_{14}^* \rho_{25} + \frac{i}{2} \Omega_{14} \rho_{34} + \frac{i}{2} \Omega_{14}^* \rho_{35} + \frac{i}{2} \Omega_{14} \rho_{45} + \frac{i}{2} \Omega_{14}^* \rho_{46},
\]

\[
\dot{\rho}_{15} = -\gamma_{15} \rho_{15} - \frac{i}{2} \Omega_{14} \rho_{14} + \frac{i}{2} \Omega_{14}^* \rho_{16} - \frac{i}{2} \Omega_{15} \rho_{25} - \frac{i}{2} \Omega_{15}^* \rho_{26} - \frac{i}{2} \Omega_{15} \rho_{35} - \frac{i}{2} \Omega_{15}^* \rho_{36} - \frac{i}{2} \Omega_{15} \rho_{45} - \frac{i}{2} \Omega_{15}^* \rho_{46}.
\]

(5)

The decay rate of the states \( |1\rangle, |2\rangle, |3\rangle, |4\rangle \) and \( |5\rangle \) are given by \( \gamma_{11}, \gamma_{12}, \gamma_{13}, \gamma_{14}, \gamma_{15} \) respectively. Again within the same weak probe limit, the coherences \( \rho_{23} \approx \rho_{34} \approx \rho_{45} \approx 0 \). Under the steady state condition i.e. \( \rho_{0j} = 0 \), the equations (5) become

\[
\dot{\rho}_{12} = \frac{i}{2} \Omega_{12} \rho_{13} + \frac{i}{2} \Omega_{12}^* \rho_{12},
\]

\[
\dot{\rho}_{13} = \frac{i}{2} \Omega_{23} \rho_{23} + \frac{i}{2} \Omega_{34} \rho_{13} + \frac{i}{2} \Omega_{13} \rho_{15} + \frac{i}{2} \Omega_{13}^* \rho_{16},
\]

\[
\dot{\rho}_{14} = \frac{i}{2} \Omega_{34} \rho_{14} + \frac{i}{2} \Omega_{14} \rho_{24} + \frac{i}{2} \Omega_{14}^* \rho_{16},
\]

\[
\dot{\rho}_{15} = \frac{i}{2} \Omega_{45} \rho_{15} + \frac{i}{2} \Omega_{15} \rho_{35} + \frac{i}{2} \Omega_{15}^* \rho_{36} + \frac{i}{2} \Omega_{15} \rho_{45} + \frac{i}{2} \Omega_{15}^* \rho_{46}.
\]

(6)

We obtain following analytical expression for \( \rho_{12} \) by solving above linear algebraic equations

\[
\rho_{12} = \frac{i}{2} \Omega_{12} \rho_{13} + \frac{i}{2} \Omega_{12}^* \rho_{12},
\]

(7)

where \( \phi = \phi_{35} + \phi_{14} + \phi_{34}. \) In order to ensure the correctness of the above approximation, in figure (2) we plot the normalized absorption of the probe field (Im(\( \rho_{12} \)T2/\( \Omega_{12} \))) as a function of normalized probe detuning (\( \delta_{12}/T_{2} \)) obtain from analytical expression of \( \rho_{12} \) as given in equation (7) as well as
the complete numerical solution of density matrix equation as stated in equation (4). The numerical solution of the density matrix equation is calculated for the parameters mentioned in the caption of figure 2. We solve 25 coupled differential equations for \( p_i \) where \( i = 1 \) to 5, \( j = 1 \) to 5, for the time 100/\( \Gamma_2 \), which is much higher than steady state time (1/\( \Gamma_2 \)). It is clear from figure 2 that the complete numerical results are in an excellent agreement with the approximated analytical solutions.

3. Results and discussions

3.1. Lineshape of the probe absorption

In this section, we study the effect of the MW field on the lineshape of the probe absorption. We consider all the control and MW fields are on the resonance. First we discuss the role of individual control fields one by one and finally the MW field. There is a broad Lorentzian profile having a full width at half maximum of \( \Gamma_2 \) which is \( 2\pi \times 6 \) MHz in this case [9, 10]. At the line center of the Lorentzian profile, the probe response develops a sharp EIT dip in the presence of the control field \( \Omega_{23} \) [9, 38]. The linewidth of the EIT dip depends on decoherence rates and the control field intensity \( \Omega_{23} \) and can be expressed as \( \Delta \omega \propto |\Omega_{23}|^2/4\Gamma_2 \) [52]. Here we have taken \( \Gamma_1 = \Gamma_3 = 0 \), as states |1⟩ and |3⟩ are the ground states. The control laser \( \Omega_{34} \) recovers the absorption against EIT i.e. EIT and absorption which is known as EITA and has been discussed in detail in our previous work [38]. When \( |\Omega_{34}| < |\Omega_{23}| \) is the linewidth of the EITA peak it is comparable with the linewidth of the EIT dip and it is broadening with the increase of the \( \Omega_{34} \). Finally when \( |\Omega_{34}| > |\Omega_{23}| \) is the linewidth of the EITA peak it will be \( \Gamma_2 \). In the presence of the control laser \( \Omega_{45} \) there will be transparency against EITA, known as EITAT, i.e. EIT, absorption and transparency, and this has also been discussed in our previous work [38]. The linewidth of the EITAT is given by the expression,

\[
\Delta \omega \propto \frac{1}{4\Gamma_2} \left[ |\Omega_{23}|^2 |\Omega_{45}| + \frac{1}{|\Omega_{23}|^2} + \frac{1}{|\Omega_{45}|^2} + \frac{|\Omega_{23}|^2 |\Omega_{45}|}{|\Omega_{23}|^2 |\Omega_{45}|} \right]^{-1}
\]

The linewidth of the EITAT is modulated by the three control field intensities \( \Omega_{23}, \Omega_{34}, \Omega_{45} \) and the decay rate of the excited states [2, 4]. Here we have taken \( \Gamma_3 = 0 \), as state |5⟩ is the metastable ground state. The MW field, \( \Omega^{\text{MW}}_{45} \) splits or shifts the EITAT dip depending upon the phase \( \phi \) as shown in figure 3.

Figure 3 shows the plot for normalized absorption (Im(\( p_{12} \))\( |p_{12}|^2 \)) vs probe detuning in the unit of \( \Gamma_2 \) for various combinations of the phase and the MW field strength. In figure 3(a) for \( \phi = 0 \) there is shift of the EITAT dip by approximately \( |\Omega^{\text{MW}}_{45}|^2/2 \) to the right while it shifts to the left by same amount for \( \phi = \pi \). For the \( \phi = \pi/2 \) there is a broadening and reduction of the EITAT dip. However with the relatively high value of the MW field, \( |\Omega^{\text{MW}}_{45}| \) \( = 2\pi \times 0.3 \) MHz there is further reduction of the EITAT dip and splitting increases and is visible as shown in figure 3(b).

3.2. Effect of the MW power

The variation of the normalized probe absorption as a function of the MW Rabi frequency, \( |\Omega^{\text{MW}}_{45}| \) in unit of \( \Gamma_2 \) is shown in figure 4. As in the presence of the MW field the EITAT dip shifts, thus with increase of the MW Rabi frequency the absorption increases and saturates to 1. The slope of the absorption profile depends upon the EITAT linewidth, and the narrower the linewidth, the sharper the slope. As we see in figure 4 the red solid curve shows the sharper rise in the absorption with the MW Rabi frequency as compared to the dashed green curve. This is because the EITAT dip linewidth is power broadened by the control laser with a higher Rabi frequency. The red solid curve saturates faster as compared to the dotted blue curve, this is because with \( \phi = \pi/2 \) the EITAT dip splits instead of shifting and hence will have a lower slope as compared to the \( \phi = 0 \) with the same Rabi frequency of the control lasers.

4. Closed loop \( \Lambda \)- versus M-system

In this section, we discuss how the constraints of a closed loop \( \Lambda \)-system [41] can be overcome by considering an M-system in which the MW field couples the unpopulated ground states. In figure 5, we study the probe absorption lineshape as a function of detuning under three-photon
resonance condition i.e. $\delta_{12} - \delta_{23} - \delta_{13} = 0$ for a closed loop $\Lambda$-system and an $M$-system. The lineshape under the approximation of no population transfer [41] is denoted by open circles in figure 5(a). A Lorentzian fit of the same (dashed blue curve) gives a linewidth of 0.028 $\Gamma_2$. A full numerical solution for this system displays a broader linewidth (0.1$\Gamma_2$) as shown by the open squares. The broadening of the lineshape is due to the population transfer to the excited states which leads to fluorescence. The solid green curve in figure 5(b) shows the full numerical solution of the density matrix equation (4) at a steady state condition for the $M$-system. It is found that the Lorentzian linewidth of the absorption spectrum is 0.014 $\Gamma_2$ as shown by the dashed blue curve. Therefore the Lorentzian linewidth of the $M$-system is much smaller than the Lorentzian linewidth of the closed loop $\Lambda$-system. Hence the population transfer can be suppressed in the former system by the MW field which is not possible in the latter. Nevertheless, the transparency window can be shifted by changing the total phase of the system at a fixed linewidth. Thus $M$-systems driven by a MW field can efficiently select the carrier frequency of the probe field.

5. Conclusions

In conclusion, we have studied how the strength and phase of the MW field can be efficiently used to manipulate the absorption property of the probe field in a five-level $M$-system. The transparency window of the probe field shifts or splits depending upon the phase of the MW field for the resonant control optical fields. The transparency window splits by an amount equal to the Rabi frequency of the MW field for $\pi/2$ phase while it shifts to higher or lower value by an amount equal to half of the Rabi frequency of the MW field for the phase 0 or $\pi$ of the same. We have also shown that $M$-system offers a linewidth an order of magnitude narrower than the probe absorption spectrum as compared to the routinely studied $\Lambda$-system even at a moderate strength MW field. This is due to the absence of MW field induced population transfer in the $M$-system in contrast to the $\Lambda$-system. The changes in the position of the transparency window in frequency space can enable us to select the desired carrier probe field frequency. Hence a narrow probe absorption spectrum opens up a new avenue for atom-based phase-dependent MW magnetometry which has important applications in MW engineering and communications.

Acknowledgments

NSM would like to acknowledge financial support from MHRD, Government of India.

References

[1] Sedlacek J A, Schwettmann A, Kübler H and Shaffer J P 2013 Phys. Rev. Lett. 111 063001
[2] Sedlacek J A, Schwettmann A, Kuhler H, Low R, Pfau T and Shaffer J P 2012 Nat. Phys. 8 819
[3] Horsley A, Du G X, Pellaton M, Affolderbach C, Mileti G and Treutlein P 2013 Phys. Rev. A 88 063407
[4] Horsley A, Du G X and Treatlein P 2015 New J. Phys. 17 112002
[5] Jiao Y, Yang Z, Li J, Raithel G, Zhao J and Jia S 2016 Phys. Rev. A 94 033823
[6] Fleischhauer M, Imamoglu A and Marangos J P 2005 Rev. Mod. Phys. 77 633–73
[7] Peng B, Ozdemir S K, Chen W, Nori F and Yang L 2014 Nature 5 5082
[8] Mishina O S, Scherman M, Lombardi P, Ortalo J, Felinto D, Sheremet A S, Bramati A, Kapriyanov D V, Laurat J and Giacobino E 2011 Phys. Rev. A 83 033809
[9] Iftiquar S M, Karve G R and Natarajan V 2008 Phys. Rev. A 77 053807
[10] Iftiquar S M and Natarajan V 2009 Phys. Rev. A 79 013808
[11] Budker D, Kimball D F, Rochester S M and Yashchuk V V 1999 Phys. Rev. Lett. 83 1767–70
[12] Wang H, Goorskey D and Xiao M 2001 Phys. Rev. Lett. 87 073601
[13] Phillips D F, Fleischhauer A, Mair A, Walsworth R L and Lukin M D 2001 Phys. Rev. Lett. 86 783–6
[14] Dey T N and Agarwal G S 2003 Phys. Rev. A 67 033813
[15] Menon S and Agarwal G S 1999 Phys. Rev. A 61 013807
[16] Lazoudis A, Kirova T, Ahmed E H, Qi P, Huennekens J and Lyrya A M 2011 Phys. Rev. A 83 063419
[17] Zhu C, Tan C and Huang G 2013 Phys. Rev. A. 87 043813
[18] Krishna A, Pandey K, Wasan A and Natarajan V 2005 Europhys. Lett. 72 221–7
[19] Pritchard J D, Maxwell D, Gauguet A, Weatherill K J, Jones M P A and Adams C 2010 Phys. Rev. Lett. 105 193603
[20] Gorskov A V, Otterbach J, Fleischhauer M, Pohl T and Lukin M D 2011 Phys. Rev. Lett. 107 133601
[21] Gorniaczyk H, Tresp C, Schmidt J, Fedder H and Hofferberth S 2014 Phys. Rev. Lett. 113 053601
[22] Noh H R and Moon H S 2015 Phys. Rev. A 92 013807
[23] Jin S, Li Y and Xiao M 1995 Opt. Commun. 119 90–6
[24] Das S, Grankin A, Iakoupov I, Brion E, Borregaard J, Boddeda R, Usmani I, Ourjoumtsev A, Grangier P and Sørensen A S 2016 Phys. Rev. A 93 040303
[25] de Echaniz S R, Greentree A D, Durrant D A V, Segal D M, Marangos J P and Vaccaro J A 2001 Phys. Rev. A 64 013812
[26] Han D, Guo H, Bai Y and Sun H 2005 Phys. Lett. A 334 243–8
[27] Bason M G, Mohapatra A K, Weatherill K J and Adams C S 2009 J. Phys. B: At. Mol. Opt. Phys. 42 075303
[28] Goren C, Wilson-Gordon A D, Rosenbluh M and Friedmann H 2004 Phys. Rev. A 69 053818
[29] Chen Y, Wei X G and Ham B S 2009 J. Phys. B: At. Mol. Opt. Phys. 42 065506
[30] Chana S R, Singh A K, Brun B, Pandey K and Natarajan V 2011 Opt. Commun. 284 4957–60
[31] Tian X D, Liu Y M, Yan X B, Cui C L and Zhang Y 2015 Opt. Commun. 345 6–12
[32] Yan D, Liu Y M, Bao Q Q, Fu C B and Wu J H 2012 Phys. Rev. A 86 023828
[33] Hong T, Cramer C, Nagourney W and Fortson E N 2005 Phys. Rev. Lett. 94 050801
[34] Pandey K, Kaundilliya D and Natarajan V 2011 Opt. Commun. 284 252–5
[35] Kumar S, Lauprêtre T, Ghosh R, Bretenaker F and Goldfarb F 2011 Phys. Rev. A 84 023811
[36] Kumar S, Lauprêtre T, Bretenaker F, Goldfarb F and Ghosh R 2013 Phys. Rev. A 88 023852
[37] Hu Y X, Miniaturo C, Wilkowski D and Grémaud B 2014 Phys. Rev. A 90 023601
[38] Pandey K 2013 Phys. Rev. A 87 043838
[39] Khazali M, Heshami K and Simon C 2015 Phys. Rev. A 91 030301
[40] Agarwal G S, Dey T N and Menon S 2001 Phys. Rev. A 64 053809
[41] Li H, Sautenkov V A, Rostovtsev Y V,潍格 G R, Hemmer P R and Scully M O 2009 Phys. Rev. A 80 023820
[42] Manjappa M, Undurti S S, Karigowda A, Narayanan A and Sanders B C 2014 Phys. Rev. A 90 043859
[43] Rajitha K V, Das S, Dey T N and Jha P K 2015 Opt. Lett. 40 2229–32
[44] Eilam A, Wilson-Gordon A D and Friedmann H 2009 Opt. Lett. 34 1834–6
[45] Kosachiov D V, Matsiov B G and Rozbodstvensky Y V 1992 J. Phys. B: At. Mol. Opt. Phys. 25 2473
[46] Radwell N, Clark T W, Piccirillo B, Barnett S M and Franke-Arnold S 2015 Phys. Rev. Lett. 114 123603
[47] Rajitha K V and Dey T N 2016 Phys. Rev. A 94 053851
[48] Lukin M D 2003 Rev. Mod. Phys. 75 457–72
[49] Hakumiyana G, Leroy C, Mirzoyan R, Pashayan-Leroy Y and Sarkisyan D 2012 Eur. Phys. J. D 66 119
[50] Sargsyan A, Tonoyan A, Hakumiyana G, Leroy C, Pashayan-Leroy Y and Sarkisyan D 2015 Europhys. Lett. 110 23001
[51] Scully M O and Zubairy M S 1997 Quantum Optics (Cambridge: Cambridge University Press) (https://doi.org/10.1017/CBO9780511813993)
[52] Fleischhauer M and Lukin M D 2002 Phys. Rev. A 65 022314