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Resonantly enhanced continuous-wave four-wave mixing with spontaneously generated coherence in a five-state cold atomic medium*

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Using Schrödinger–Maxwell formalism, we propose and analyse a continuous-wave four-wave mixing (FWM) scheme for the generation of coherent light in a five-state double-Λ atomic system with or without spontaneously generated coherence (SGC) based on electromagnetically induced transparency (EIT). We derive the corresponding explicit analytical expressions for the generated FWM field under the steady-state condition. The influence of hyperfine sublevel and SGC effect on the amplitude of the generated FWM field is predicted in detail via the derived analytical expressions. We also give a brief discussion on the experimental realization of the proposed scheme.

Keywords: continuous wave, four-wave mixing, EIT
PACC: 4250, 4265M

1. Introduction

Electromagnetically induced transparency (EIT) is the optical transparency of a three-level medium at a resonant transition induced by application of a coherent electromagnetic field at an adjacent transition. The accompanying reduction of the group velocity of light and large enhancement of the nonlinearity of the coherent atomic medium at the few-photon level by many orders of magnitude in the EIT transparency windows has been both intensively and actively investigated in connection with many potential applications, especially low-intensity nonlinear optics (i.e., it can reach a level of few photons), optical multi-wave mixing, ultraslow optical bright and dark solitons based on Raman excitation, ultraslow photon entanglement, quantum information processing and quantum information storage and so on.[1−22] In a recent experiment, Yan *et al* demonstrated that electromagnetically induced transparency in a standard Λ-type configuration could be used to suppress both single-photon and two-photon absorptions simultaneously.[3] Harris *et al* proposed the use of EIT to suppress absorption of the short-wavelength light generated in a four-wave mixing (FWM) scheme and showed that the FWM efficiency could be greatly enhanced.[5] Recently, Deng *et al* reported the first theoretical investigation of optical coherent four-wave mixing (OCFWM) with a weak probe wave based on electromagnetically induced transparency and showed that such a scheme could lead to many orders of magnitude of enhancement in the amplitude of the generated wave in a typical four-level atomic system.[9] Later on, Wu *et al* analysed and discussed a four-wave-mixing (FWM) scheme in a five-level atomic system and hyper-Raman scattering (HRS) in resonant coherent media by the use of EIT, which led to suppressing both single-photon, two-photon and three-photon absorptions in both FWM and HRS schemes and enabling the four-wave mixing to proceed through real, resonant intermediate states without absorption loss.[11,12] The recent studies show that when the generated field becomes sufficiently intense in such an efficient four-wave mixing process, however, it can be

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absorbed dramatically. Consequently, a robust three-photon destructive interference between the two different excitation channels occurs, resulting in a saturated FWM production. Quite recently, a new scheme proposed by Deng and Payne retains the significantly enhanced conversion efficiency enabled by ultraslow propagation of pump waves, yet is also capable of inhibiting and delaying the onset of detriment three-photon destructive interference which limits the further growth of FWM.[14,15] In this paper, we propose and analyse a continuous-wave (CW) four-wave mixing scheme via resonant nonlinear optical process with or without spontaneously generated coherence (SGC) effect in a five-level double-Λ atomic system modified from the generic four-level double-Λ system. We present the corresponding explicit analytical expressions for the generated FWM field under the steady-state condition. The influence of hyperfine sublevel and SGC effect on the amplitude of the generated FWM field is predicted. To conclude we give a brief discussion on the experimental realization.

This paper is organized as follows. In Section 2, the theoretical model is described and the Schrödinger–Maxwell equations of motion that govern the dynamics of such a five-level system are solved analytically in the presence of the SGC effect. In Section 3, we analyse and discuss the influence of hyperfine sublevel and SGC effect on the amplitude of the generated FWM field. A summary appears in Section 4.

2. FWM scheme and solutions of the Schrödinger–Maxwell equations of motion

Let us consider the FWM scheme depicted in Fig.1. Three laser fields (frequencies $\omega_{\rm p}$, $\omega_{\rm c1}$ and $\omega_{\rm c2}$, respectively) induce the FWM processes $|1\rangle \rightarrow |4\rangle(|3\rangle) \rightarrow |2\rangle \rightarrow |5\rangle \rightarrow |1\rangle$ and efficiently generate the coherent radiation field (amplitude being $E_{\rm m}(\omega_{\rm m}, z)$) at a frequency $\omega_{\rm m} = \omega_{\rm p} - \omega_{\rm c1} + \omega_{\rm c2}$. Because of energy conservation there is overall four-photon resonance. At the same time, a couple of two-photon resonance with the $|1\rangle \leftrightarrow |2\rangle$ transition is also satisfied.

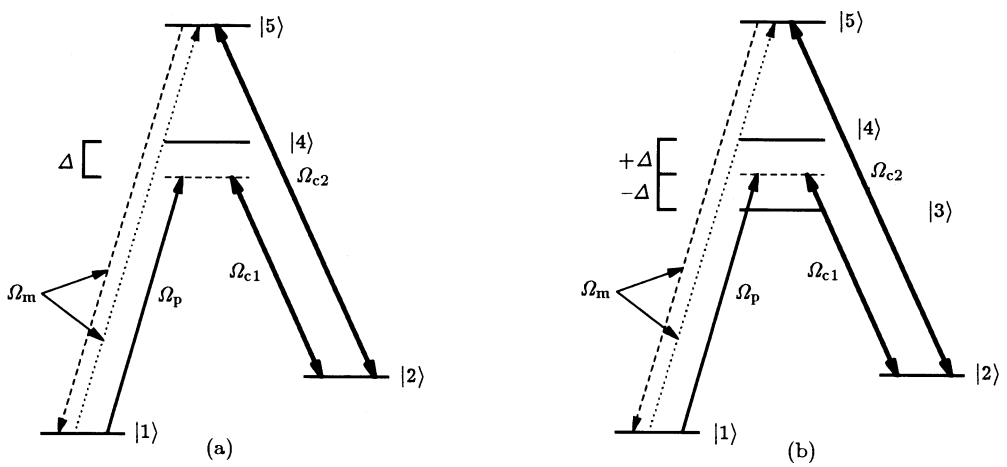


Fig.1. (a) Four-wave mixing scheme in a generic double-Λ system. (b) Four-wave mixing scheme in a modified double-Λ system including a hyperfine doublet structure labelled $|3\rangle$ and $|4\rangle$ with a frequency difference of 2Δ. In this case the lifetime-broadened atomic system is driven by a weak CW pump laser field (frequency $\omega_{\rm p}$ and Rabi frequency $\Omega_{\rm p}$) and two strong CW control laser fields (frequencies $\omega_{\rm c1}$ and $\omega_{\rm c2}$ and Rabi frequencies $\Omega_{\rm c1}$ and $\Omega_{\rm c2}$) to generate a coherent FWM field (frequency $\omega_{\rm m}$ and Rabi frequency $\Omega_{\rm m}$).

Using the electro-dipole and rotating-wave approximations, the semiclassical Hamiltonian describing the atom–field interaction for the system under study can be written in the Schrödinger picture as

$$\begin{aligned} H =&\sum_{j=1}^{5} E_j^{(a)}|j\rangle\langle j| - \hbar(g\Omega_{\rm p}{\rm e}^{{\rm i}\theta_{\rm p}}|3\rangle\langle 1| + {\rm h.c.}) \\ &- \hbar(\Omega_{\rm p}{\rm e}^{{\rm i}\theta_{\rm p}}|4\rangle\langle 1| + {\rm h.c.}) \\ &- \hbar(f\Omega_{\rm c1}{\rm e}^{{\rm i}\theta_{\rm c1}}|3\rangle\langle 2| + {\rm h.c.}) \\ &- \hbar(\Omega_{\rm c1}{\rm e}^{{\rm i}\theta_{\rm c1}}|4\rangle\langle 2| + {\rm h.c.}) \\ &- \hbar(\Omega_{\rm c2}{\rm e}^{{\rm i}\theta_{\rm c2}}|5\rangle\langle 2| + {\rm h.c.}) \\ &- \hbar(\Omega_{\rm m}{\rm e}^{{\rm i}\theta_{\rm m}}|5\rangle\langle 1| + {\rm h.c.}), \end{aligned} \tag{1}$$

where h.c. means Hermitian conjugation. The quantities $\theta_n = \boldsymbol{k}_n \cdot \boldsymbol{r} - \omega_n t$ (n=p, m, c1, c2) are the phase factors of the positive frequency optical fields. Ω_n (n=p, c1, c2, m) denote one-half Rabi frequencies for the relevant laser-driven transitions, i.e. $\Omega_{\rm p} = \mu_{41}E_{\rm p}/(2\hbar)$, $\Omega_{\rm c1} = \mu_{42}E_{\rm c1}/(2\hbar)$, $\Omega_{\rm c2} = \mu_{52}E_{\rm c2}/(2\hbar)$, $\Omega_{\rm m} = \mu_{51}E_{\rm m}/(2\hbar)$, and $\mu_{ij} = \boldsymbol{\mu}_{ij} \cdot \hat{e}_{\rm L}$ with $\hat{e}_{\rm L}$ being the unit polarization vector of the laser fields. The parameters $g = \mu_{31}/\mu_{41}$ and $f = \mu_{32}/\mu_{42}$ are the ratios of the dipole moments, with μ_{ij} standing for the dipole moment elements for the transition between levels $|i\rangle$ and $|j\rangle$. $E_j^{(a)} = \hbar\omega_j (j = 1-5)$ are the energy of the atomic state $|j\rangle$. For simplicity, in what follows we take $E_1^{(a)} = 0$ for the ground state $|1\rangle$ as the energy origin. Turning to the interaction picture, the Hamiltonian can be rewritten as (taking $\hbar = 1$)

$$\begin{aligned} H_0 =&(\omega_{\rm p} - \omega_{\rm c1})|2\rangle\langle 2| + \omega_{\rm p}|3\rangle\langle 3| \\ &+ \omega_{\rm p}|4\rangle\langle 4| + (\omega_{\rm p} - \omega_{\rm c1} + \omega_{\rm c2})|5\rangle\langle 5|, \\ H_{\rm int} =&(\Delta_{\rm p} - \Delta_{\rm c1})|2\rangle\langle 2| + (\Delta_{\rm p} - \Delta)|3\rangle\langle 3| \\ &+ (\Delta_{\rm p} + \Delta)|4\rangle\langle 4| + (\Delta_{\rm p} - \Delta_{\rm c1} + \Delta_{\rm m})|5\rangle\langle 5| \\ &- (g\Omega_{\rm p}{\rm e}^{{\rm i}\vec{k}_{\rm p}\cdot\vec{r}}|3\rangle\langle 1| + \Omega_{\rm p}{\rm e}^{{\rm i}\vec{k}_{\rm p}\cdot\vec{r}}|4\rangle\langle 1| \\ &+ f\Omega_{\rm c1}{\rm e}^{{\rm i}\vec{k}_{\rm c1}\cdot\vec{r}}|3\rangle\langle 2| + \Omega_{\rm c1}{\rm e}^{{\rm i}\vec{k}_{\rm c1}\cdot\vec{r}}|4\rangle\langle 2| \\ &+ \Omega_{\rm c2}{\rm e}^{{\rm i}\vec{k}_{\rm c2}\cdot\vec{r}}|5\rangle\langle 2| + \Omega_{\rm m}{\rm e}^{{\rm i}\vec{k}_{\rm m}\cdot\vec{r}}|5\rangle\langle 1| + {\rm h.c.}), \end{aligned} \tag{2}$$

where $\Delta_{\rm p} = (\omega_4 + \omega_3)/2 - \omega_{\rm p}$, $\Delta_{\rm c1} = (\omega_4 + \omega_3)/2 - \omega_2 - \omega_{\rm c1}$, and $\Delta_{\rm m} = \omega_5 - \omega_2 - \omega_{\rm c2}$ are the detunings of the laser-driven fields from the corresponding transitions. $\Delta = (\omega_4 - \omega_3)/2$ is a half frequency difference of two closely lying hyperfine sublevels.

Let us assume that the wavefunction has the form

$$\begin{aligned} |\Psi\rangle =&A_1|1\rangle + A_2{\rm e}^{{\rm i}(\vec{k}_{\rm p}-\vec{k}_{\rm c1})\cdot\vec{r}}|2\rangle \\ &+ A_3{\rm e}^{{\rm i}\vec{k}_{\rm p}\cdot\vec{r}}|3\rangle + A_4{\rm e}^{{\rm i}\vec{k}_{\rm p}\cdot\vec{r}}|4\rangle \\ &+ A_5{\rm e}^{{\rm i}(\vec{k}_{\rm p}-\vec{k}_{\rm c1}+\vec{k}_{\rm c2})\cdot\vec{r}}|5\rangle, \end{aligned} \tag{3}$$

where $A_j (j = 1-5)$ stand for the time-dependent probability amplitudes of the atomic states.

Making use of the Schrödinger equation in the interaction picture ${\rm i}\partial_t|\Psi\rangle = H_{\rm int}|\Psi\rangle$ and the Weisskopf–Wigner theory, the equations of motion for the slowly varying probability amplitude of the atomic wavefunctions can be readily obtained as follows:

$$(\partial_t + \tilde{\gamma}_2)A_2 = {\rm i}f^*\Omega_{\rm c1}^*A_3 + {\rm i}\Omega_{\rm c1}^*A_4 + {\rm i}\Omega_{\rm c2}^*A_5, \tag{4a}$$

$$\begin{aligned} (\partial_t + \tilde{\gamma}_3)A_3 =&{\rm i}g\Omega_{\rm p}A_1 + {\rm i}f\Omega_{\rm c1}A_2 - \frac{1}{2}(\sqrt{\gamma_{41}\gamma_{31}}\cos\theta_1 \\ &+ \sqrt{\gamma_{42}\gamma_{32}}\cos\theta_2)A_4, \end{aligned} \tag{4b}$$

$$\begin{aligned} (\partial_t + \tilde{\gamma}_4)A_4 =&{\rm i}\Omega_{\rm p}A_1 + {\rm i}\Omega_{\rm c1}A_2 - \frac{1}{2}(\sqrt{\gamma_{41}\gamma_{31}}\cos\theta_1 \\ &+ \sqrt{\gamma_{42}\gamma_{32}}\cos\theta_2)A_3, \end{aligned} \tag{4c}$$

$$(\partial_t + \tilde{\gamma}_5)A_5 = {\rm i}\Omega_{\rm c2}A_2 + {\rm i}\Omega_{\rm m}{\rm e}^{-{\rm i}\delta\vec{k}\cdot\vec{r}}A_1, \tag{4d}$$

where we have introduced the new definitions $\tilde{\gamma}_2 = \gamma_2 + {\rm i}(\Delta_{\rm p} - \Delta_{\rm c1}) = \gamma_{21} + {\rm i}(\Delta_{\rm p} - \Delta_{\rm c1})$, $\tilde{\gamma}_3 = \gamma_3 + {\rm i}(\Delta_{\rm p} - \Delta) = \gamma_{31} + \gamma_{32} + {\rm i}(\Delta_{\rm p} - \Delta)$, $\tilde{\gamma}_4 = \gamma_4 + {\rm i}(\Delta_{\rm p} + \Delta) = \gamma_{41} + \gamma_{42} + {\rm i}(\Delta_{\rm p} + \Delta)$ and $\tilde{\gamma}_5 = \gamma_5 + {\rm i}(\Delta_{\rm p} - \Delta_{\rm c1} + \Delta_{\rm m}) = \gamma_{51} + \gamma_{52} + {\rm i}(\Delta_{\rm p} - \Delta_{\rm c1} + \Delta_{\rm m})$. The notation $\delta\boldsymbol{k} = \boldsymbol{k}_{\rm p} - \boldsymbol{k}_{\rm c1} + \boldsymbol{k}_{\rm c2} - \boldsymbol{k}_{\rm m}$ represents the phase matching condition and $\gamma_{ij} (i, j = 1-5)$ are the decay rates from states $|{\rm i}\rangle$ to $|j\rangle$. The terms including $\frac{1}{2}\sqrt{\gamma_{41}\gamma_{31}}\cos\theta_1 A_{3(4)}\left(\frac{1}{2}\sqrt{\gamma_{42}\gamma_{32}}\cos\theta_2 A_{3(4)}\right)$ represent the quantum interference effect resulting from the cross coupling pathways between spontaneous emissions $|4\rangle \to |1\rangle(|4\rangle \to |2\rangle)$ and $|3\rangle \to |1\rangle(|3\rangle \to |2\rangle)$, i.e., the spontaneously generated coherence (SGC).[19] $\theta_1(\theta_2)$ is the angle between the two induced dipole moments $\mu_{41}(\mu_{42})$ and $\mu_{31}(\mu_{32})$. If the above angles subject to $\theta_1 = \theta_2 = 0$ (parallel dipole moments), the system exhibits maximum SGC effect, while if the angles $\theta_1 = \theta_2 = \pi/2$ (orthogonal dipole moments), no SGC effect occurs. Alternatively, it is worthwhile to point out that only for small energy spacing between the two higher levels $|4\rangle$ and $|3\rangle$ is the SGC effect remarkable; as for large energy spacing, the oscillatory terms will average out to zero and such an effect will vanish. In the following, we consider all laser fields to be continuous and analyse the FWM process in the steady state. Following the standard method described before in Refs.[4, 12], the pump laser field $\Omega_{\rm p}$ are assumed to be sufficiently weak ($|\Omega_{\rm p}| \ll |\Omega_{\rm c1,c2}|$), so that in essence all of the atomic population remains in the ground state $|1\rangle$: i.e., using the nondepleted

ground-state approximation $A_1 \simeq 1$. Therefore, we can then obtain straightforwardly the steady-state solutions of Eqs.(4) and the slowly varying component of the atomic polarization of the generated FWM field: i.e., $P_{\rm m} = N\mu_{15}A_5A_1^*$ with N being the atomic number density in the medium. Without loss of generality, we assume that the driving and FWM generated fields propagate in the z direction, hereafter we take $\delta\boldsymbol{k} = \delta k \cdot \hat{z}$. Based on Eqs.(4a)–(4d) and setting $\partial_t A_j = 0$ for the steady-state solution, we can derive the probability amplitude of the atomic state $|5\rangle$ as follows:

$$A_5 = -\mathrm{i}L + \mathrm{i}R\Omega_{\rm m}\mathrm{e}^{-\mathrm{i}\delta kz}, \tag{5}$$

with

$$L = \frac{(\eta f^* - f^* g\tilde{\gamma_4} + \eta g - \tilde{\gamma_3})\Omega_{\rm p}\,\Omega_{\rm c1}^*\,\Omega_{\rm c2}}{\tilde{\gamma_2}(\eta^2 - \tilde{\gamma_3}\tilde{\gamma_4})\tilde{\gamma_5} + (\eta f^* + \eta f - |f|^2\tilde{\gamma_4} - \tilde{\gamma_3})\tilde{\gamma_5}|\Omega_{\rm c1}|^2 + (\eta^2 - \tilde{\gamma_3}\tilde{\gamma_4})|\Omega_{\rm c2}|^2},$$

$$R = \frac{\tilde{\gamma_2}(\eta^2 - \tilde{\gamma_3}\tilde{\gamma_4}) + (\eta f^* + \eta f - |f|^2\tilde{\gamma_4} - \tilde{\gamma_3})|\Omega_{\rm c1}|^2}{\tilde{\gamma_2}(\eta^2 - \tilde{\gamma_3}\tilde{\gamma_4})\tilde{\gamma_5} + (\eta f^* + \eta f - |f|^2\tilde{\gamma_4} - \tilde{\gamma_3})\tilde{\gamma_5}|\Omega_{\rm c1}|^2 + (\eta^2 - \tilde{\gamma_3}\tilde{\gamma_4})|\Omega_{\rm c2}|^2},$$

$$\eta = \frac{1}{2}(\sqrt{\gamma_{41}\gamma_{31}}\cos\theta_1 + \sqrt{\gamma_{42}\gamma_{32}}\cos\theta_2).$$

In order to describe correctly the propagation of the generated FWM field, the atomic equations of motion must be simultaneously solved with Maxwell's equation in a self-consistent manner. In the limit of plane waves and slowly varying amplitude approximations, the amplitude of the generated FWM field $E_{\rm m} = E_{\rm m}(\omega_{\rm m}, z)$ obeys Maxwell's equation

$$(\partial_z + 1/V_g\,\partial_t)E_{\rm m} = \mathrm{i}\omega_{\rm m}P_{\rm m}/(2c\varepsilon_0), \tag{6}$$

where $V_{\rm g}$ are the group velocity of the generated FWM field. c and ε_0 are the light speed and permittivity of free space, respectively.

Making full use of $P_{\rm m} = N\mu_{15}A_5A_1^*$, $\Omega_{\rm m} = \mu_{51}E_{\rm m}/(2\hbar)$ and Eq.(5) in Eq.(6), we have

$$(\partial_t + V_{\rm g}\partial_z)\Omega_{\rm m} = V_{\rm g}\cdot M(L - R\Omega_{\rm m}\mathrm{e}^{-\mathrm{i}\delta kz}), \tag{7}$$

where $M = \omega_{\rm m}N|\mu_{15}|^2/(4\hbar\varepsilon_0 c)$ is the propagation constant, which is related to the frequently used oscillator strength of the transition $|1\leftrightarrow|5\rangle$.

Neglecting pump depletion, under phase matching condition $\delta k \approx 0$ we can readily solve Eq.(7) for the slowly varying field amplitude $\Omega_{\rm m}(\omega_{\rm m}, z)$. Applying the boundary condition $\Omega_{\rm m}|_{z=0} = 0$ yields explicit analytical expressions of the generated FWM field

$$\Omega_{\rm m}(\omega_{\rm m}, z) = \frac{L}{R}(1 - \mathrm{e}^{-MRz}). \tag{8}$$

3. Analysis and discussion for the generated FWM field

In the following analysis, for the CW control and pump laser fields we will only consider the case of resonant excitations, i.e. $\Delta_{\rm p,c1,m} = 0$, in which the EIT effects are the most prominent. Under the condition of perfect phase matching $\delta k \approx 0$, in this section we have analysed the influence of the ratios of dipole moments g and f on the amplitude of the generated FWM field $\Omega_{\rm m}$ without SGC effect, as shown in Fig.2. Note that, in this paper, parameters $\Omega_{\rm p,c1,c2,m}$, Δ, $\gamma_{ij}(i,j = 1-5)$ and M are scaled with γ_4. From Fig. 2, we find that the amplitude of the generated FWM field increases monotonically with the increasing of the propagation distance z firstly, then the amplitude of the FWM field reaches the saturation value and is independent of the propagation distance z for the five-level atomic system. This result is the behaviour of a quantum destructive interference between the two different excitation channels: $|5\rightarrow|1\rangle$ (coupling excitation channel, the dashed line in Fig. 1) and $|1\rightarrow|5\rangle$ (back-coupling excitation channel, the dotted line in Fig. 1).[14, 15] That is to say, after a characteristic propagation distance when the generated FWM field has been sufficiently intense, an efficient back-coupling excitation channel to the FWM generated state $|5\rangle$ becomes important, which is π out of phase with respect to the coupling excitation channel. This back-coupling channel simultaneously gives rise to a destructive interference pathway, thereby suppresses the amplitude of the generated FWM state $|5\rangle$. As a matter of fact, this is a consequence of the competition between the creation of quantum coherence by the coupling excitation channel and the destruction of this coherence by the back-coupling excitation channel. From Fig.2, it is worth pointing out that when the ratios of the dipole moments $g = 1, f = -1$

or $g = -1, f = 1$, the amplitude of the generated FWM field is greatly enhanced. This is due to the fact that in the four-level system ($f = g = 0$), the finite detuning Δ is associated with an ac-Stark effects, which leads to intensity-dependent dynamical phase shifts of the fields. These phase shifts have a detrimental influence, leading to impaired phase matching and inefficient frequency conversion. However, these nonlinear phase shifts in such a five-level scheme can be eliminated by ac-Stark compensation only when we choose atomic states such that the dipole moment for one of the four transitions $|1\rangle \to |3\rangle$, $|1\rangle \to |4\rangle$, $|2\rangle \to |3\rangle$, and $|2\rangle \to |4\rangle$ has a different sign to the other three, i.e., $g = -f$. This can easily be accomplished by using different hyperfine sublevels.[16] In Fig.3, we plot the amplitude of the generated FWM field $\Omega_{\rm m}(\omega_{\rm m}, z)$ as a function of propagation distance z with the SGC effect. We find that in the presence of the SGC effect, the amplitude of the generated FWM field clearly decreases for the case of $g = -f$. On the basis of the above analysis, it is obvious that in order to achieve large amplitude of the generated FWM field, we should eliminate the existence of the SGC effect. This can be realized by using the orthogonal dipole moments.

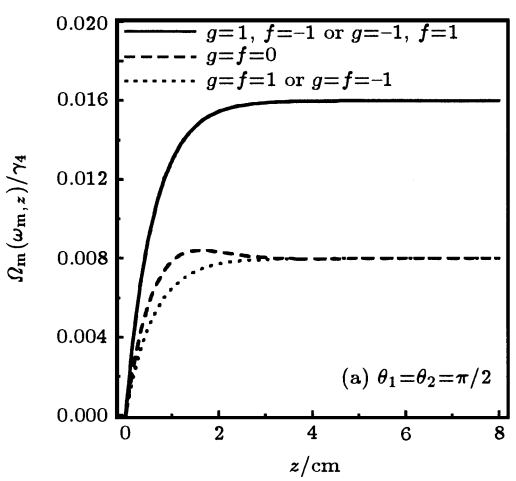


Fig.2. Plots of the amplitude of the generated FWM field $|\Omega_{\rm m}(\omega_{\rm m}, z)|$ as a function of propagation distance z without SGC effect. The fitting parameters: $\Omega_{\rm p} = 0.02\gamma_4$, $\Omega_{\rm c1} = 0.5\gamma_4$, $\Omega_{\rm c2} = 0.2\gamma_4$, $\gamma_2 = 0$, $\gamma_{41} = \gamma_{42} = \gamma_{31} = \gamma_{32} = 0.5\gamma_4$, $\gamma_5 = 0.2\gamma_4$, $\Delta = 2\gamma_4$, $\Delta_{\rm p,c1,m} = 0$, $\dfrac{\omega_{\rm m} N|\mu_{15}|^2}{4\hbar\varepsilon_0 c} = \gamma_4 {\rm cm}^{-1}$.

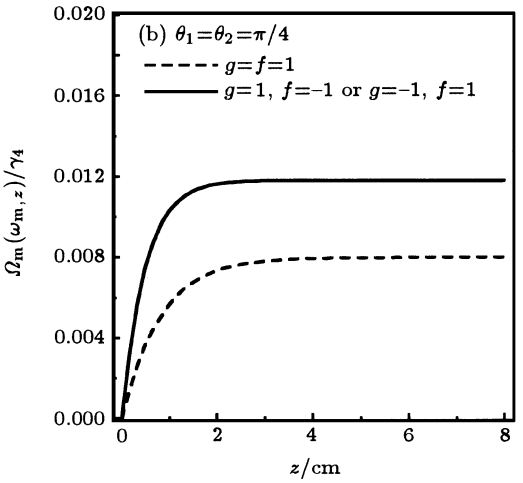


Fig.3. Plots of the amplitude of the generated FWM field $|\Omega_{\rm m}(\omega_{\rm m}, z)|$ as a function of propagation distance z with SGC effect. The fitting parameters: $\Omega_{\rm p} = 0.02\gamma_4$, $\Omega_{\rm c1} = 0.5\gamma_4$, $\Omega_{\rm c2} = 0.2\gamma_4$, $\gamma_2 = 0$, $\gamma_{41} = \gamma_{42} = \gamma_{31} = \gamma_{32} = 0.5\gamma_4$, $\gamma_5 = 0.2\gamma_4$, $\Delta = 2\gamma_4$, $\Delta_{\rm p,c1,m} = 0$, $\dfrac{\omega_{\rm m} N|\mu_{15}|^2}{4\hbar\varepsilon_0 c} = \gamma_4 {\rm cm}^{-1}$.

Now we give a brief discussion on the experimental realization of our scheme. An experimental candidate for the proposed system can be found in ^{87}Rb (nuclear spin $I = 3/2$) cold atoms with the designated states chosen as follows: Levels $|1\rangle$ and $|2\rangle$ can be, for example, the $M_F = \pm 1$ Zeeman sublevels of the $5S_{1/2}(F = 1)$ ground state of cold rubidium. Levels $|3\rangle$ and $|4\rangle$ are the $M_F = 0$ Zeeman sublevels of $5P_{1/2}(F = 1)$ and $5P_{1/2}(F = 2)$, and $5P_{3/2}(F = 1, M_F = 0)$ as $|5\rangle$. The optimum detuning is exactly midway between the hyperfine states $|3\rangle$ and $|4\rangle$ as shown in Fig. 1(b), leading to the ratios of the dipole moments $|f| = |g| = 1$ and $f = -g$.[16]

4. Conclusions

In summary, using the Schrödinger–Maxwell equations of the five-level atomic system and the light fields we have demonstrated a CW four-wave mixing scheme in the presence of the SGC effect in highly resonant media. The proposed scheme may be useful, in principle, for generation of high-frequency and/or high power laser light.

Acknowledgments

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