Large enhancement of four-wave mixing by suppression of photon absorption from electromagnetically induced transparency

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(Received 26 July 2002; revised manuscript received 7 August 2002; published 28 January 2003)

We analyze a four-wave-mixing (FWM) scheme in a five-level atomic system based on electromagnetically induced transparency (EIT). We show that EIT suppresses both two-photon and three-photon absorptions in the FWM scheme and enables the four-wave mixing to proceed through real, resonant intermediate states without absorption loss. The scheme results in a several orders of magnitude increase in the FWM efficiency in comparison with a recent scheme [Phys. Rev. Lett. **88**, 143902 (2002)] and may be used for generating short-wavelength radiation at low pump intensities.

DOI: 10.1103/PhysRevA.67.013811 PACS number(s): 42.50.Gy, 42.65.Ky, 42.50.Hz

Electromagnetically induced transparency (EIT) modifies both absorptive and dispersive properties of an absorbing medium, which may lead to suppression of lower-order susceptibilities and enhancement of higher-order susceptibilities [1–8]. Because of the near-resonant atom-photon interaction and slow group velocity [9–11], nonlinear optical phenomena may be observed at low light intensities approaching single-photon energy levels [12]. There are many theoretical and experimental studies in the literature dealing with efficiency generation of short-wavelength coherent light based on EIT and/or coherent population trapping [2,3,13–15]. In particular, it is of interest to explore EIT and atomic coherence in applications of vacuum ultraviolet (VUV) radiation generation from processes of multiphoton conversion of long-wave photons. Harris *et al.* proposed 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 can be greatly enhanced [16]. Agarwal *et al.* proposed an efficient, degenerate FWM scheme based on amplification of the pump field in a coherently driven system by a strong off-resonant control field [17]. Sokolov *et al.* reported efficient generation of short-wavelength light via a phase-coherent Raman process [18]. Recently, Deng *et al.* proposed a FWM scheme based on EIT and associated slow light propagation [19]. Their calculations showed that a many orders of magnitude increase in the FWM efficiency may be obtained.

EIT usually refers to suppression of single-photon absorption by destructive quantum interference in an absorbing medium. However, under suitable conditions, single-photon as well as multiphoton absorptions may be suppressed simultaneously. Agarwal *et al.* showed that in a four-level Ladder-type system, two-photon absorption can be selectively suppressed or enhanced [20]. This was confirmed experimentally by Gao *et al.* [21]. In a recent experiment, Yan *et al.* demonstrated that EIT in a standard Λ-type configuration can be used to suppress both single-photon and two-photon absorptions simultaneously [22]. Here we show that suppression of the two-photon absorption can be effectively used to enhance the FWM process in a scheme [Fig. 1(a)] modified from the one proposed by Deng *et al.* [Fig. 1(b)]. Both schemes rely on EIT and the phase matching manifested by the slow group velocity to dramatically enhance the FWM efficiency as analyzed by Deng *et al.* In the FWM scheme [Fig. 1(b)] from Deng *et al.*, the intermediate state is virtual and thus the effective two-photon Rabi frequency usually is very small. The modified scheme [Fig. 1(a)] uses the real intermediate state $|3\rangle$ and two separate pump fields ω_1 and ω_2 that are resonantly coupled with the intermediate state $|3\rangle$. The EIT suppresses both the two-photon absorption $|0\rangle \to |2\rangle \to |3\rangle$ and three-photon absorption $|0\rangle \to |2\rangle \to |3\rangle \to |4\rangle$, enabling the FWM process to proceed through the real intermediate states and to be immune to the absorption loss. Because of large enhancement of the FWM process through the stepwise two-photon resonance and interference cancellation of the resonant nonlinear absorptions in the modified scheme, the FWM efficiency can be further increased by several orders of magnitude more. Such improvement in the FWM efficiency makes the scheme suitable for generating short-wavelength radiation at low pump intensities and may be applicable in other applications.

Consider the FWM scheme depicted in Fig. 1(a). Three pump fields (frequencies ω_p, ω_1, and ω_2, respectively) in-

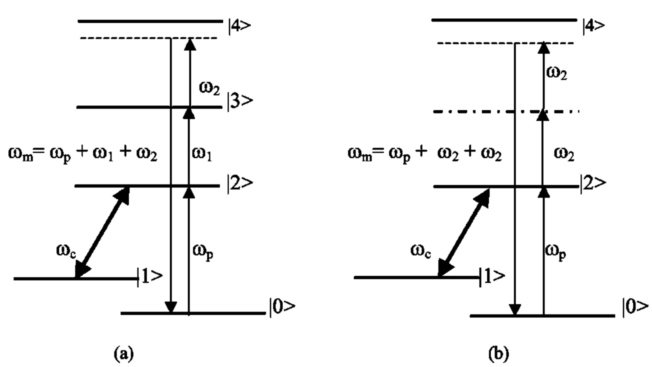

FIG. 1. (a) Modified FWM scheme with a real intermediate state $|3\rangle$ coupled by two separate laser fields ω_1 and ω_2. (b) FWM scheme of Ref. [19] that results in a large enhancement of the FWM efficiency manifested by EIT and associated slow group velocities.

duce the FWM process $|0\rangle\rightarrow|2\rangle\rightarrow|3\rangle\rightarrow|4\rangle\rightarrow|0\rangle$ and generate the radiation field at a frequency $\omega_m=\omega_p+\omega_1+\omega_2$. A control field with frequency ω_c couples states $|1\rangle$ and $|2\rangle$ and creates EIT. When ω_1 and ω_2 satisfy the two-photon resonance ($|0\rangle\rightarrow|2\rangle\rightarrow|3\rangle$) and the three-photon resonance ($|0\rangle\rightarrow|2\rangle\rightarrow|3\rangle\rightarrow|4\rangle$) conditions, the two nonlinear absorption processes are suppressed by the EIT-induced destructive quantum interference [22,23]. As will be shown below, the absorption suppression and the enhancement of the FWM through the resonant intermediate state $|3\rangle$ lead to an orders of magnitude increase of the FWM efficiency in comparison with the FWM scheme of Ref. [19] [Fig. 1(b)].

The Hamiltonian describing the system shown in Fig. 1(a) is given by ($\hbar=1$)

$$\mathcal{H}=\sum_{j=1}^{4}\epsilon_j|j\rangle\langle j|-(\Omega_c e^{-i\theta_c}|2\rangle\langle 1|+\text{H.c.}),$$
$$-(\Omega_p e^{-i\theta_p}|2\rangle\langle 0|+\text{H.c.})-(\Omega_1 e^{-i\theta_1}|3\rangle\langle 2|+\text{H.c.}),$$
$$-(\Omega_2 e^{-i\theta_2}|4\rangle\langle 3|+\text{H.c.})-(\Omega_m e^{-i\theta_m}|4\rangle\langle 0|+\text{H.c.}), \tag{1}$$

where $\theta_j=\omega_j t-\vec{k}_j\cdot\vec{r}$, Ω_j ($j=1,2,c,p,m$) denote one-half of the Rabi frequencies for the respective transitions [$\Omega_m=D_{04}E(\omega_m)/(2\hbar)$, $\Omega_c=D_{12}E(\omega_c)/(2\hbar)$, $\Omega_p=D_{02}E(\omega_p)/(2\hbar)$, $\Omega_1=D_{23}E(\omega_1)/(2\hbar)$, and $\Omega_2=D_{34}E(\omega_2)/(2\hbar)$, with D_{ln} denoting the dipole moment for the transition between levels $|l\rangle$ and $|n\rangle$], and ϵ_j is the energy of the atomic state $|j\rangle$ (taking $\epsilon_0=0$ for the ground state $|0\rangle$). In the interaction picture, the Hamiltonian can be written as

$$H=-\Delta\omega_c|1\rangle\langle 1|-\Delta\omega_p|2\rangle\langle 2|-\Delta\omega_3|3\rangle\langle 3|-\Delta\omega_4|4\rangle\langle 4|$$
$$-(\Omega_c e^{i\vec{k}_c\cdot\vec{r}}|2\rangle\langle 1|+\Omega_p e^{i\vec{k}_p\cdot\vec{r}}|2\rangle\langle 0|+\Omega_1 e^{i\vec{k}_1\cdot\vec{r}}|3\rangle\langle 2|$$
$$+\Omega_2 e^{i\vec{k}_2\cdot\vec{r}}|4\rangle\langle 3|+\Omega_m e^{i\vec{k}_m\cdot\vec{r}}|4\rangle\langle 0|+\text{H.c.}), \tag{2}$$

where $\Delta\omega_p=\omega_p-\epsilon_2$ is the single-photon detuning, $\Delta\omega_c=(\omega_p-\omega_c)-\epsilon_1$ and $\Delta\omega_3=(\omega_p+\omega_1)-\epsilon_3$ are two separate two-photon detunings, respectively, and $\Delta\omega_4=(\omega_p+\omega_1+\omega_2)-\epsilon_4\equiv\Delta\omega_3+\Delta\omega_{43}$ [$\Delta\omega_{43}=\omega_2-(\epsilon_4-\epsilon_3)$] represents the three-photon detuning (denoted as δ_m in Ref. [19]).

From the Schrödinger equation in the interaction picture, $i(\partial|\Psi\rangle/\partial t)=H|\Psi\rangle$, and defining the atomic state as $|\Psi\rangle=A_0|0\rangle+A_1 e^{i(\vec{k}_p-\vec{k}_c)\cdot\vec{r}}|1\rangle+A_2 e^{i\vec{k}_p\cdot\vec{r}}|2\rangle+A_3 e^{i(\vec{k}_p+\vec{k}_1)\cdot\vec{r}}|3\rangle+A_4 e^{i\vec{k}_m\cdot\vec{r}}|4\rangle$, one readily obtains

$$\frac{\partial}{\partial t}A_0=i\Omega_p^*A_2+i\Omega_m^{\star}A_4, \tag{3a}$$

$$\left(\frac{\partial}{\partial t}+\gamma_1-i\Delta\omega_c\right)A_1=i\Omega_c^*A_2, \tag{3b}$$

$$\left(\frac{\partial}{\partial t}+\gamma_2-i\Delta\omega_p\right)A_2=i\Omega_pA_0+i\Omega_cA_1+i\Omega_1^*A_3, \tag{3c}$$

$$\left(\frac{\partial}{\partial t}+\gamma_3-i\Delta\omega_3\right)A_3=i\Omega_1A_2+i\Omega_2^*e^{i\delta\vec{k}\cdot\vec{r}}A_4, \tag{3d}$$

$$\left(\frac{\partial}{\partial t}+\gamma_4-i\Delta\omega_4\right)A_4=i\Omega_2e^{-i\delta\vec{k}\cdot\vec{r}}A_3+i\Omega_mA_0, \tag{3e}$$

where $\delta\vec{k}=\vec{k}_m-(\vec{k}_p+\vec{k}_1+\vec{k}_2)$ and γ_j ($j=1-4$) are the decay rates. As in Ref. [19], we consider the coupling field ω_c and pump fields ω_1 and ω_2 to be continuous, but the pump field ω_p and the generated field ω_m have a time-dependent beam profile. Following the standard procedure described in Ref. [19], we take $A_0\approx 1$ (assuming $\Omega_c\gg\Omega_p$) and carry out Fourier transformations $A_j(t)\sim\int\alpha_j(\omega)\exp(-i\omega t)d\omega$ ($j=1,2,3,$ and 4) for Eqs. (3b)–(3e) to obtain the following equations:

$$\Delta\tilde{\omega}_c\alpha_1+\Omega_c^*\alpha_2=0, \tag{4a}$$

$$\Omega_c\alpha_1+\Delta\tilde{\omega}_p\alpha_2+\Omega_1^*\alpha_3=-W_p, \tag{4b}$$

$$\Omega_1\alpha_2+\Delta\tilde{\omega}_3\alpha_3+\Omega_2^*e^{i\delta\vec{k}\cdot\vec{r}}\alpha_4=0, \tag{4c}$$

$$\Omega_2e^{-i\delta\vec{k}\cdot\vec{r}}\alpha_3+\Delta\tilde{\omega}_4\alpha_4=-W_m, \tag{4d}$$

where $\Delta\tilde{\omega}_c\equiv\omega+\Delta\omega_c+i\gamma_1$, $\Delta\tilde{\omega}_p\equiv\omega+\Delta\omega_p+i\gamma_2$, and $\Delta\tilde{\omega}_j\equiv\omega+\Delta\omega_j+i\gamma_j$ ($j=3,4$), and W_p and W_m are the Fourier transformations of Ω_p and Ω_m, respectively.

It is then straightforward, though tedious, to solve Eq. (4) and obtain the following results:

$$\alpha_2=-\frac{\Delta\tilde{\omega}_c\Omega_1^*\Omega_2^*e^{i\delta\vec{k}\cdot\vec{r}}}{S}W_m-\frac{(\Delta\tilde{\omega}_3\Delta\tilde{\omega}_4-|\Omega_2|^2)\Delta\tilde{\omega}_c}{S}W_p, \tag{5a}$$

$$\alpha_4=-\left[\frac{1}{\Delta\tilde{\omega}_4}+\frac{|\Omega_2|^2(\Delta\tilde{\omega}_p\Delta\tilde{\omega}_c-|\Omega_c|^2)}{S\Delta\tilde{\omega}_4}\right]W_m$$
$$-\frac{\Delta\tilde{\omega}_c\Omega_1\Omega_2e^{-i\delta\vec{k}\cdot\vec{r}}}{S}W_p, \tag{5b}$$

where $S=(\Delta\tilde{\omega}_p\Delta\tilde{\omega}_c-|\Omega_c|^2)(\Delta\tilde{\omega}_3\Delta\tilde{\omega}_4-|\Omega_2|^2)-|\Omega_1|^2\Delta\tilde{\omega}_c\Delta\tilde{\omega}_4$.

The pump field W_p and the generated FWM field W_m obey Maxwell's equations,

$$\frac{\partial W_p}{\partial z}-i\frac{\omega}{c}W_p=i\kappa_{02}\alpha_2, \tag{6a}$$

$$\frac{\partial W_m}{\partial z}-i\frac{\omega}{c}W_m=i\kappa_{04}\alpha_4, \tag{6b}$$

where $\kappa_{02(04)}=2N\omega_{p(m)}|D_{02(04)}|^2/(c\hbar)$, with N denoting atomic concentration. Letting $\kappa_{02}\alpha_2=g_2e^{i\delta k\cdot z}W_m+f_2W_p$ and $\kappa_{04}\alpha_4=g_4e^{-i\delta k\cdot z}W_p+f_4W_m$, Eq. (6) can be written as

$$\frac{\partial W_p}{\partial z}-ib_2W_p=ig_2\tilde{W}_m, \tag{7a}$$

$$\frac{\partial \widetilde{W}_m}{\partial z} - i b_4 \widetilde{W}_m = i g_4 W_p, \tag{7b}$$

where $b_2 = \omega/c + f_2$, $b_4 = \omega/c + f_4 + \delta k$, and $\widetilde{W}_m = W_m e^{i\delta k \cdot z}$. The solutions to Eq. (7) are

$$W_p = Q_2 e^{ik_+ z} + R_2 e^{ik_- z}, \tag{8a}$$

$$W_m = (Q_4 e^{ik_+ z} + R_4 e^{ik_- z}) e^{-i\delta k \cdot z}, \tag{8b}$$

where the constants Q_j and R_j are determined by the boundary conditions at $z=0$, and

$$k_{\pm} = \frac{\omega}{c} + \frac{f_2 + f_4 + \delta k}{2} \pm \sqrt{\left(\frac{f_2 - f_4 - \delta k}{2}\right)^2 + g_2 g_4}, \tag{9}$$

where

$$f_2 = -\kappa_{02} \frac{(\Delta\widetilde{\omega}_3 \Delta\widetilde{\omega}_4 - |\Omega_2|^2)\Delta\widetilde{\omega}_c}{S}, \quad g_2 = -\kappa_{02} \frac{\Delta\widetilde{\omega}_c \Omega_1^* \Omega_2^*}{S}, \tag{10}$$

$$f_4 = -\kappa_{04}\left[\frac{1}{\Delta\widetilde{\omega}_4} + \frac{|\Omega_2|^2(\Delta\widetilde{\omega}_p \Delta\widetilde{\omega}_c - |\Omega_c|^2)}{S\Delta\widetilde{\omega}_4}\right],$$

$$g_4 = -\kappa_{04} \frac{\Omega_1 \Omega_2 \Delta\widetilde{\omega}_c}{S}. \tag{11}$$

For the FWM emission process, $W_m(z=0)=0$. The generated FWM field W_m is then given by

$$W_m = W_p(z=0)\,\frac{g_4}{k_+ - k_-}(e^{ik_+ z} - e^{ik_- z}) e^{-i\delta k \cdot z}. \tag{12}$$

The power of the generated FWM field is $P_m = \hbar \omega_m A |\Omega_m|^2/(2\gamma_4 \sigma_{04})$, where Ω_m is the Fourier transformation of W_m in Eq. (12), A is the area, and σ_{04} is the atomic absorption cross section. The efficiency of the FWM process can be defined as $(\hbar=1)$ [12]

$$\epsilon = \frac{(1/\omega_m)\int P_m dt}{(1/\omega_p)\int P_p dt}. \tag{13}$$

Equation (7b) shows that the driving source of the generated four-wave-mixing field W_m is $ig_4 W_p$. Equation (12) gives the generated FWM field W_m at arbitrary frequency detunings $\Delta\omega_j$ ($j=p$, c, 3, or 4), the Fourier transform of which can be used to calculate the generated FWM power. The FWM efficiency can then be derived from Eq. (13).

As discussed in Ref. [19], a phase-matched FWM field can be produced if the group velocity of the probe wave ω_p is substantially reduced by EIT and about equal to the group velocity of the generated FWM wave ω_m. In order to show that the FWM scheme is capable of generating the FWM field at low pump intensities, we choose the pump field Rabi frequencies Ω_1, Ω_2, and Ω_p well below the saturation levels and present numerical calculations in Fig. 2. Figure 2(a) plots the generated FWM field W_m versus the two-photon detuning $\Delta\omega_3$ and Fig. 2(b) plots the FWM field W_m versus the detuning $\Delta\omega_{43}$. The calculations show that W_m is maximized at $\Delta\omega_3=0$ and $\Delta\omega_{43}=0$, demonstrating the enhancement of the FWM via the resonant intermediate state manifested by the EIT-induced suppression of the two-photon and three-photon absorptions.

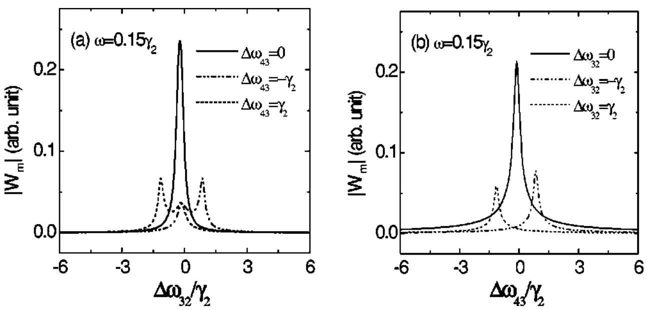

FIG. 2. Generated FWM field W_m versus (a) $\Delta\omega_3$ and (b) $\Delta\omega_{43}$. The parameter values are $\Omega_c=\gamma_2$, $\Omega_p=0.05\gamma_2$, $\Omega_1=\Omega_2=0.1\gamma_2$, $\gamma_3=0.2\gamma_2$, $\gamma_4=0.1\gamma_2$, $\gamma_1=0.001\gamma_3$, $\Delta\omega_c=\Delta\omega_p=0$, $\omega=0.15\gamma_2$, and $\kappa_{02}=100\kappa_{04}=\gamma_2/z$. Similar results are obtained for other ω values. The FWM W_m is maximized at the two-photon resonance $\Delta\omega_3=0$ and the three-photon resonance $\Delta\omega_4=0$, demonstrating enhancement of the FWM efficiency from the EIT-induced suppression of the two-photon and three-photon absorptions.

The EIT width (half-width at half maximum) is given by $\Gamma = \sqrt{\gamma_2^2 + 4\Omega_c^2} - \gamma_2$. For the light pulse of the pump field ω_p to acquire the slow group velocity and minimum absorption loss, the pump pulse width (in time) should be $\tau \geqslant 1/\Gamma$. When this condition is satisfied, the adiabatic approximation is valid [6]. Under such conditions and assuming a Gaussian pulse shape $[\Omega_p(t)=\Omega_{p0}\exp(-t^2/\tau^2)]$ for the pump field ω_p, we present the calculated FWM efficiency from Eq. (13) in Fig. 3 for several sets of parameters. As expected from Fig. 2, the FWM efficiency is optimized near the two-photon and three-photon resonances with the maximum value occurring at $\Delta\omega_p=\Delta\omega_{32}=\Delta\omega_{43}=0$, where the EIT suppression of the linear and nonlinear absorption leads to the minimized absorption loss for the three pump fields and results in a greatly enhanced FWM efficiency.

For a quantitative comparison with Ref. [19] that treated the EIT-enhanced FWM process without the resonant intermediate state [Fig. 1(b)], we consider the FWM process under the resonant atom-field interactions $\Delta\omega_j=0$ ($j=p,c$) in which EIT-meditated slow group velocity for the pump field ω_p enhances the FWM process and the generated FWM wave also travels with an ultraslow group velocity. We further assume $\Omega_c \gg \Omega_p, \Omega_1$, and Ω_2, i.e., considering the FWM at low pump intensities. Under these conditions, the generated FWM field reduces to

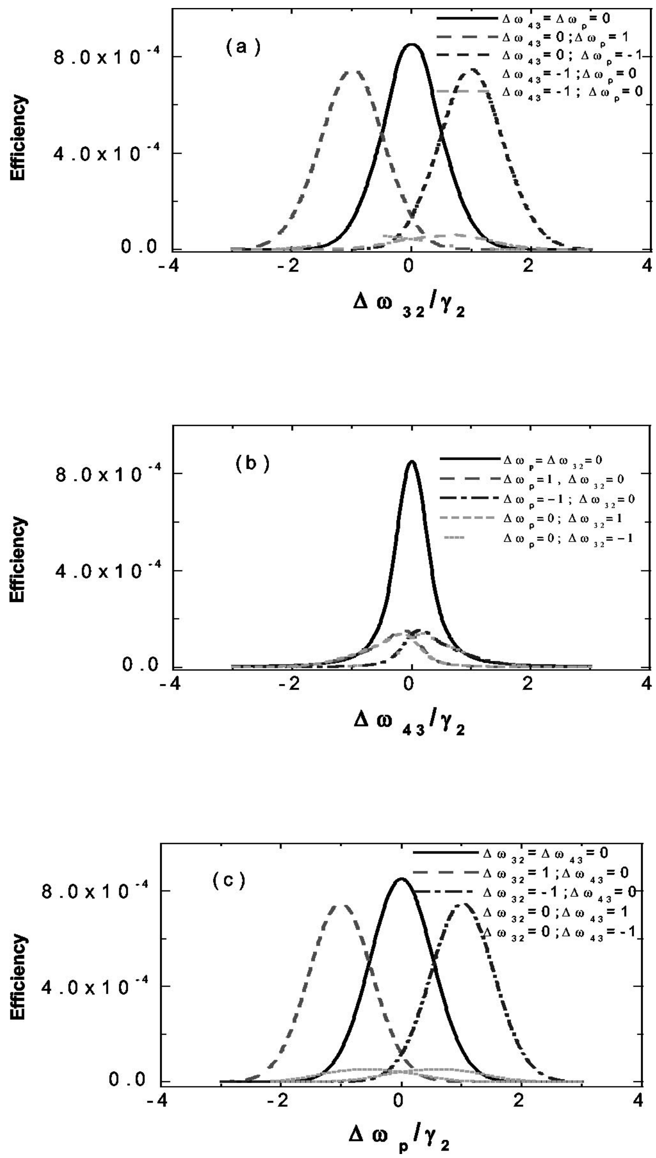

FIG. 3. Calculated FWM efficiency ϵ versus (a) $\Delta\omega_{32}$, (b) $\Delta\omega_{43}$, and (c) $\Delta\omega_p$. The parameter values are $\Omega_c=0.5\gamma_2$, $\Omega_p=0.02\gamma_2$, $\Omega_1=\Omega_2=0.1\gamma_2$, $\gamma_3=0.14\gamma_2$, $\gamma_4=0.015\gamma_2$, $\gamma_1=0.001\gamma_3$, $\Delta\omega_c=0$, and $\kappa_{02}=100\kappa_{04}=\gamma_2/z$. All the detunings labeled within the figure are measured in γ_2. The FWM W_m is maximized at the two-photon resonance $\Delta\omega_3=0$ and the three-photon resonance $\Delta\omega_4=0$. The maximum efficiency is $\sim 10^{-3}$ for the chosen parameters.

$$W_m \approx W_m' \frac{\Omega_1\Omega_2}{\Omega_{32}^{(2)}\Delta\tilde{\omega}_3}, \tag{14}$$

where W_m' [denoted as W_{m30} in Eq. (4) of Ref. [19]] is the generated FWM field for the model system of Ref. [19]. Ω_1 (Ω_2) is one-half of the single-photon Rabi frequency for the transition $|2\rangle \to |3\rangle$ ($|3\rangle \to |4\rangle$) [Fig. 1(a)] while $\Omega_{32}^{(2)}$ is one-half of the Rabi frequency for the two-photon process via a virtual state considered in Ref. [19] and usually is very small due to the large frequency detuning Δ' from the real intermediate state ($\Omega_{32}^{(2)} \approx \Omega_1\Omega_2/\Delta'$). Under the two-photon resonant condition $\Delta\tilde{\omega}_3 = i\gamma_3$ [$\Delta\omega_3=(\omega_p+\omega_1)-\epsilon_3=0$], which leads to the maximum FWM efficiency as shown before, one has

$$\left|\frac{W_m}{W_m'}\right| \approx \left|\frac{\Omega_1\Omega_2}{\Omega_{32}^{(2)}\Delta\tilde{\omega}_3}\right| \approx \left|\frac{\Delta'}{\gamma_3}\right| \gg 1. \tag{15}$$

As an example, taking $\gamma_3 \sim 10^6$ Hz and $\Delta' \geq 10^9$ Hz (usually $\gg 10^9$ Hz for nonresonant two-photon transitions), the ratio is $\geq 10^3$. Therefore, taking advantage of the EIT suppression of the two-photon absorption, one can use the resonant pump fields in the FWM process and obtain orders of magnitude increases in the FWM efficiency in addition to the great enhancement of the FWM process based on the single-photon EIT and slow group velocity enhancement analyzed in Ref. [19]. An experimental candidate for the proposed system can be found in ^{85}Rb atoms with the designated states chosen as follows: $5S_{1/2}$, $F=2$ as $|0\rangle$, $5S_{1/2}$, $F=3$ as $|1\rangle$, $5P_{1/2}$ as $|2\rangle$, $5D_{3/2}$ as $|3\rangle$, and $nP_{3/2}$ as $|4\rangle$ ($n>10$). The respective transitions are $|0\rangle \to |2\rangle$ and $|1\rangle \to |2\rangle$ at 795 nm ($\gamma_2 \approx 5.9$ MHz), $|2\rangle \to |3\rangle$ at 762 nm ($\gamma_3 \approx 0.8$ MHz), and $|3\rangle \to |4\rangle$ at 1.3–1.5 μm ($\gamma_4 \approx 0.09$ MHz), all accessible with diode lasers. The wavelength of the generated Raman light is ~ 300 nm. With the atomic density $N \sim 10^{12}$ cm^{-3}, $L=2$ mm (for the Rb atoms in a magneto-optical trap where the Doppler broadening is negligible), $\Omega_c=\gamma_2$, $\Omega_p=0.05\gamma_2$, and $\Omega_1=\Omega_2=0.1\gamma_2$, the calculated efficiency from Eq. (13) is $\epsilon \approx 10^{-3}$. Without EIT, the dominant, weak-field single-photon absorption will deplete the first pump field completely in a distance $\ll L$, which reduces the estimated Raman efficiency (assuming the comparable Raman gain with or without EIT) to $<10^{-7}$. Then, the expected EIT enhancement is $>10^4$.

In conclusion, we have analyzed an EIT-enhanced FWM scheme and shown that the FWM field can be efficiently generated with weak pump fields. Our scheme uses the resonant intermediate states and one more laser field to facilitate the resonant two-photon and three-photon couplings. Due to the EIT-induced suppression of the two-photon and three-photon absorptions, an additional increase of the FWM efficiency by several orders of magnitude can be obtained. Thus, the FWM scheme based on EIT and EIT-induced suppression of the nonlinear absorptions is capable of generating the FWM field with extremely high efficiency and may be used to produce a short wavelength, coherent radiation at low pump intensities in applications such as nonlinear laser spectroscopy and quantum nonlinear optics.

Y.W. is supported by the National Fundamental Research Program of China (Grant No. 2001CB309310), the National Natural Science Foundation of China through Grant Nos. 90108026, 60078023, and 10125419, and by the Chinese Academy of Sciences through the 100 Talents Project and Grant No. KJCX2-W7-4. Y.Z. is supported by the National Science Foundation (Grant Nos. 0140032 and 9729133) and the Office of Naval Research (Grant No. N00014-01-1-0754).

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