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Phase-Conjugate Amplification Through Transverse Optical Zeeman Pumping in Resonant Doppler-Broadened Degenerate Four-Wave Mixing

C. Schmidt-Iglesias and G. Orriols

Departament de Física, Universitat Autònoma de Barcelona, E-08193 Bellaterra (Barcelona), Spain

F. Pi

Escola d'Optica, Universitat Politècnica de Catalunya, E-08222 Terrassa (Barcelona), Spain

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Abstract. Degenerate four-wave mixing in an inhomogeneously broadened folded three-level system is considered to analyse theoretically the generation of phase-conjugate signals in optically dense media. The counter-propagating pump waves and the copropagating probe and forward pump waves are taken to be orthogonally polarized and interact selectively with the coupled one-photon transitions such that ground-state Zeeman coherence plays the dominant role in phase-conjugate wave generation. It is shown that even in a purely absorptive resonant interaction, transverse optical pumping gives rise to amplified reflection and coupled-mode oscillation, being the required pump intensities very low compared to that needed for saturation of the optical transitions.

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Optical phase conjugation via backward degenerate four-wave mixing (DFWM) has been arousing increasing interest over the past few years [1]. In particular, a lot of attention has been paid to resonantly enhanced DFWM in gas media since its establishment as powerful spectroscopic tool [2]. In the pursuit of high phase conjugation efficiencies, resonant gas media provide the remarkable advantage of requiring low incident power levels for reasonable phase-conjugate (PC) mirror operation [3–5]. Atomic vapors close to optical resonance have also permitted the observation of amplified PC reflection [6] and thus real-time intracavity aberration correction [7] or self-oscillations [8] by incorporating such a gaseous PC mirror as one cavity end.

Amplification, oscillation, and, in general, the PC reflectivity behaviour obtained from saturated absorbing media has been usually investigated with simplified theoretical models of stationary atoms [9–12]. These theories do not describe adequately the DFWM properties related to the Doppler-broadened nature of gas media, where the effect of atomic motion leads to the washout of the optically induced population or coherence gratings [13] and severely alters the PC emission line shapes [4].

Rather than resorting to numerical analysis, Doppler effects have been considered with the restrictive assumption of a single saturating incident pump beam [14]. Only very recently, distinct physical mechanisms giving rise to DFWM signals, as the cross-population [15] or Zeeman coherence grating mechanism [16], have been selectively explored in the Doppler regime (when the Rabi frequency of the intense pump beams is smaller than the Doppler width) considering the relevant situation of two fully saturating cross-polarized pump beams. But keeping spectroscopic applications in mind, these studies have been restricted to the limit of optically thin media and thus do not master the full potential of resonant DFWM for phase conjugation.

In this paper we extend these theoretical treatments and include the coupled-mode propagation effects, essential for eventual PC amplification. In particular, we concentrate on the generation of ground-state Zeeman coherences that are known to strongly influence the PC emission line-shape characteristics [5, 12, 16–18]. Zeeman coherence may result in a very efficient nonlinear mechanism for phase conjugation and we show that, in the Doppler regime, PC reflectivities exceeding unity and oscillation are possible at pump intensities much lower than that required for appreciable saturation of the optical transition. Laser detuning from optical resonance is not required, in contradistinction to previously reported PC amplification in saturable absorbers, and thus pump-imbalanced irradiation neither prevents oscillation. The physical phenomenon responsible for it is transverse optical pumping (TOP) that drastically reduces the field absorption and is well known in the context of high-resolution spectroscopy [19].

The DFWM configuration used to convert TOP in the relevant DFWM generation mechanism is discussed in next section. Section 1 also states the assumptions involved in modeling the four-wave interaction. The basic characteristics of the induced nonlinear polarization are exposed in Sect. 2, and Sect. 3 analyses the resulting PC reflectivity behaviour.

1. Description of the DFWM Interaction

We consider a gas medium composed of Λ-shaped three-level absorbers with two optical transitions of frequencies $W_{01}=W_0+\delta$ and $W_{02}=W_0-\delta$, such that the two sublevels 1 and 2 in the ground state are separated by a variable Zeeman splitting 2δ.

In the DFWM process shown in Fig. 1, the nonlinear medium interacts with two collinear sets of counter-propagating plane waves given by

$$\mathbf{E}_i=\sum_{\substack{\mu=\pm\\ i=1,2}}\tfrac{1}{2}A_i^{\mu}\mathbf{e}_i\{\exp[-i(Wt-\mu kz)]+\text{c.c.}\}\,. \tag{1}$$

The four interacting waves, although degenerate in frequency, are discriminated on the basis of their orthogonal polarization, and the optical field $\mathbf{E}_i$ selectively couples to the $0-i$ transition with a detuning $\Delta_i=(W_{0i}-W)$ that in the case of degenerate Zeeman sublevels reduces to $\Delta_0=W_0-W$. The amplitudes A_1^- and A_2^+ describe the two pump waves that are counter-

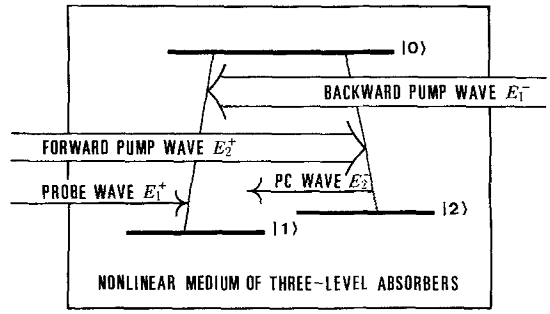

Fig. 1. Illustration of the geometry and the transition scheme of the analysed resonant DFWM interaction with orthogonally polarized pump beams $\mathbf{E}_2^+$, $\mathbf{E}_1^-$ and probe beam $\mathbf{E}_1^+$

propagating and orthogonally polarized, as required for vectorial phase conjugation [20]. A_1^+ and A_2^- denote, respectively, the amplitudes of probe wave and generated PC signal wave which travel in opposite directions due to phase-matching constraints.

The field-atom interaction gives rise to a macroscopic polarization responsible for the PC wave generation as for the mutual influence of the four optical fields, as they propagate along the medium. To simplify the propagation part of the DFWM problem, we follow a procedure similar to that adopted by Abrams and Lind [10] in the case of two-level systems. It consists of neglecting pump depletion and assuming A_1^+, $A_2^- \ll A_1^-$, A_2^+ all over the interaction length L. Introducing then the susceptibilities X_i^{μ} and K_i^{μ}, the induced nonlinear polarization may be expressed in the form

$$\begin{aligned}\mathbf{P}=\tfrac{1}{2}\varepsilon_0\{&[X_1^-A_1^-\exp(-ikz)\\&+(X_1^+A_1^+ +K_1^-A_2^{-*})\exp(ikz)]\mathbf{e}_1\\&+[X_2^+A_2^+\exp(ikz)+(X_2^-A_2^- +K_2^+A_1^{+*})\\&\times\exp(-ikz)]\mathbf{e}_2\}\exp(-iWt)+\text{c.c.}\,,\end{aligned} \tag{2}$$

and, further neglecting pump absorption effects, the amplitudes A_1^+ and A_2^- are found to satisfy, in the slowly varying envelope approximation, the set of two coupled equations:

$$\begin{aligned}&\partial A_1^+/\partial z=ik/2(X_1^+A_1^+ +K_1^-A_2^{-*}),\\&\partial A_2^{-*}/\partial z=ik/2(X_2^{-*}A_2^{-*}+K_2^{+*}A_1^+)\end{aligned} \tag{3}$$

with z-independent absorption and mixing coefficients, X_i^{μ} and K_i^{μ}. Using the boundary condition $A_2^{-*}(L)=0$, the PC reflectivity is simply given by

$$R=\frac{|A_2^{-*}(0)|^2}{|A_1^+(0)|^2}=\left|\frac{K_2^{+*}\tanh(GLk/2)}{G+X\tanh(GLk/2)}\right|^2 \tag{4}$$

with the gain parameter

$$G=(X^2-K^2)^{1/2} \tag{5}$$

and

$$X = \mathrm{i}/2(X_2^{-*} - X_1^+), \quad K = (K_1^- K_2^{+*})^{1/2}. \tag{6}$$

2. Nonlinear Susceptibility

The nonlinear polarization (2) is determined within the semiclassical approach of the density-matrix formalism for a three-level system. Following the treatments of [15, 16], the effect of the strong pump fields is dealt with exactly, whatever their intensity, while the weak probe and generated field interaction is described perturbatively up to first order.

In the theoretical approach the field-free population is assumed to be equally distributed between the two lower sublevels and obeys the standard Maxwellian velocity distribution, with most probable thermal velocity u. Equal population decay rates are supposed for the lower levels, $\gamma_1 = \gamma_2 = \gamma$, while the excited one relaxes faster so that $r = \gamma_0/\gamma \gg 1$. The coherence decay rates are allowed to include the effect of phase-perturbing collisions by assuming $\Gamma_{ij} = (\gamma_i + \gamma_j)/2 + \Gamma_{ij}^c$, and we also admit $\Gamma_{01} = \Gamma_{02} = \Gamma$.

The steady-state solution for the components of the first-order susceptibility adquires the following form in the rotating wave approximation

$$\begin{aligned} K_j^\mu = & -2i\alpha_0/k \int dkv \exp[-(v/u)^2] Z_j^\mu \\ & \times [\sqrt{\gamma\Gamma I_j} N_j^\mu - I_{3-j}\sqrt{\gamma\Gamma I_j} \\ & \times (N_{3-j}^\mu Z_{3-j}^{-\mu*} + N_j^\mu Z_j^\mu)/S_0 + I_{3-j}\sqrt{I_{3-j} I_j} \\ & \times Z_j^{-\mu}(N_{3-j}^0 Z_{3-j}^{\mu*} + N_j^0 Z_j^{-\mu})/S_0 S_1 \\ & -\sqrt{I_{3-j} I_j}(N_{3-j}^0 Z_{3-j}^{-\mu*} + N_j^0 Z_j^{-\mu})/S_0], \end{aligned} \tag{7}$$

$$\begin{aligned} X_j^\mu = & -2i\alpha_0/k \int dkv \exp[-(v/u)^2] Z_j^{-\mu} \\ & \times [N_j^0 + \sqrt{\gamma\Gamma I_j} N_j^\mu - I_{3-j}\sqrt{\gamma\Gamma I_j} \\ & \times (N_j^\mu Z_j^{-\mu} + N_{3-j}^\mu Z_{3-j}^{\mu*})/S_0 \\ & + I_{3-j} I_j Z_{3-j}^{-\mu*}(N_{3-j}^0 Z_{3-j}^{-\mu*} + N_j^0 Z_j^\mu)/S_0 S_1 \\ & - I_{3-j}(N_{3-j}^0 Z_{3-j}^{-\mu*} + N_j^0 Z_j^{-\mu})/S_0]. \end{aligned} \tag{8}$$

Here N_j^μ are the first-order Fourier components of the $N_0 - N_j$ population differences modulated at frequency $-2kz$ due to the spatial coupling between the weak $\mathbf{E}^\mu$ wave and the counter-propagating $\mathbf{E}^{-\mu}$ pump wave, N_j^0 represent the corresponding spatially uniform zero-order components induced solely by the pump waves, and are well-known in the theory of three-level saturation spectroscopy [21]. The saturation denominators S_n are given by

$$S_0 = 1 + \Gamma_{12}^c/\gamma - iD_0 + Z_j^\mu I_{3-j} + Z_{3-j}^{-\mu*} I_j$$

and

$$S_1 = 1 + \Gamma_{12}^c/\gamma - \mathrm{i}D_1 + Z_j^{-\mu} I_{3-j} + Z_{3-j}^{\mu*} I_j \tag{9}$$

with dimensionless detuning parameters

$$Z_j^\mu = [1 - \mathrm{i}(\Delta_j + \mu kv)/\Gamma]^{-1}$$

and

$$D_n = 2(\delta + nkv)\Gamma/\gamma. \tag{10}$$

The resonant unsaturated field-attenuation coefficient is

$$\alpha_0 = (k|d|^2 N/4\hbar\varepsilon_0 \Gamma\sqrt{\pi} ku)$$

being d the electric dipole moment and N the atomic density. The pump intensities $I_1 = |A_1^-|^2/I_s$ and $I_2 = |A_2^+|^2/I_s$ are normalized to the resonant saturation intensity $I_s = (4\hbar^2\gamma_0\Gamma/r|d|^2)$, where the inverse dependence on r of I_s expresses the TOP influence on the effective saturation of the three-level system.

While (7, 8) in conjunction with (4, 5) constitute a complete solution of the mixing problem, their form is rather complicated, and thus some intuition into their nature is afforded by identifying the contribution of the distinct grating mechanisms in DFWM [13]. The first term in the mixing coefficient K_j^μ (7) describes the saturated cross-population grating mechanism, the last the Zeeman coherence mechanism, whereas the second and third ones represent high-order Zeeman coupling or saturation effects relevant at high saturation regimes. At moderate saturating pump intensities (below the Rabi regime, i.e. $\sqrt{4I_j\gamma_0\Gamma/r} \ll ku$), only the last Zeeman coherence term is effective for PC wave generation due to atomic motional washout effects on the remaining contributions. Similarly, in the absorption coefficient X_j^μ (8) only the first and last terms, respectively, describing the saturated normal population and low order coupling Zeeman coherence contributions, remain significant after the velocity integration.

The last Zeeman coherence term in (7, 8) is responsible for noteworthy features of the first-order susceptibility. Figure 2 shows the numerically computed spectra of the mixing ($K_2^{+*} = -K_1^-$) and absorption ($X_1^+ = -X_2^{-*}$) susceptibility components, obtained in the Doppler limit

$$(ku = 120\gamma_0 \gg \Gamma, \ |A_j^\mu|^2/\gamma\Gamma)$$

by varying the Zeeman splitting parameter δ, after choosing $\Delta_0 = 0$, $\Gamma_{0j}^c = 0$, $r = 100$, and $I_1 = I_2 = I_p = 3$. The most important feature of Fig. 2 is the simultaneousness of a pronounced decrease in the absorption (X_1^+ in Fig. 2a) and an efficient nonlinear generation mechanism (K_2^{+*} in Fig. 2b) in a very narrow region around $\delta = 0$. The width and strength of these central resonances are closely determined by the relaxation ratio Γ_{12} of the Zeeman coherence. In particular, if dephasing collisions are supposed at high enough rate

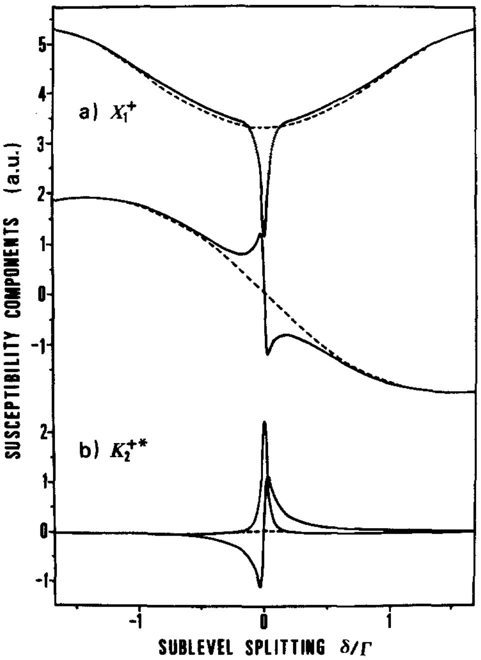

Fig. 2. (a) Absorption and (b) mixing susceptibility components as a function of the normalized sublevel splitting $2\delta/\Gamma$ after choosing $I_1=I_2=3$, $\Delta_0=0$, $r=100$, $\Gamma_{ij}^c=0$, and $ku=120\gamma_0$. The dashed curves are obtained for $\Gamma_{12}^c=\gamma_0$ and the other parameters unchanged. The imaginary parts are represented by dispersive curves and the real parts by bell-shaped ones

such that $\Gamma_{12}^c=\gamma_0$, maintaining other constants unchanged, the Zeeman coherence practically vanishes and the central dip on X_1^+ disappears as does the PC wave generation described by K_2^{+*}. The narrow central features exhibit a power broadened halfwidth, fundamentally characterized by the denominator S_0 (9) in the last Zeeman coherence term of (7), which is somewhat smaller than

$$\Delta_z=\gamma(1+\Gamma_{12}^c/\gamma+I_1+I_2) \tag{11}$$

due to the atomic velocity distribution effects. They are related to the so-called nonabsorption resonance [19] arising from a resonant mechanism of TOP that traps the atoms in a nonabsorbent coherent state of the ground sublevels, with an efficiency that increases with the relaxation ratio r. This atomic population trapping does not reduce but on the contrary enhances the generation efficiency of the resonant medium in contradistinction to previously reported TOP effects in DFWM [5, 12]. In those cases, by either neglecting atomic motion effects [12] or using a different experimental polarization scheme for the incident waves [5], the cross-population mechanism, obviously affected by transverse pumping, does also intervene in the PC wave generation reducing the line-center emission efficiency. As a result generation is most favoured in the wings ($\delta>0$), where the dispersive contribution of Zeeman coherence dominates and PC emission exhibits typical lineshape splitting. On the contrary, in our DFWM configuration, maximum mixing efficiency results at line center and furthermore the occurrence of a nonabsorption resonance permits its predominance over absorption. In fact, due to the resonant mechanism of TOP, mixing predominance establishes at very low pump intensities, provided the collisional relaxation of Zeeman-coherence is negligible and the upper state relaxes much faster than the degenerate lower state ($r\gg 1$). This is clearly shown in Fig. 3 where the line center variation of K_2^{+*} and X_1^+ with the normalized pump intensity is plotted for different values of r and Γ_{12}^c. In addition to the higher mixing efficiencies obtained at lower pump intensities (the saturation intensity scales linearly with $1/r$) when $r\geqq 100$ and $\Gamma_{12}^c=0$, mixing dominates absorption already for $I_p\approx I_s$ even when saturation on the optical transitions is negligible.

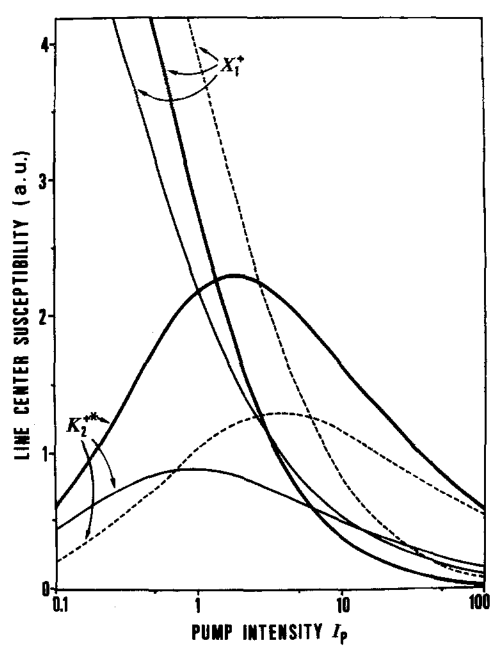

Fig. 3. Variation of the line center ($\delta=0$) absorption X_1^+ and mixing K_2^{+*} susceptibility components with the normalized pump intensity I_p (equal-intensity pump irradiance) for different values of the parameter r defined as the ratio of upper-to-lower-level relaxation rates and of the collisional Zeeman coherence decay rate Γ_{12}^c. The thick line curves correspond to $r=100$, $\Gamma_{12}^c=0$, the thin continuous line curves to $r=1$, $\Gamma_{12}^c=0$, and the broken line curves to $r=100$ and $\Gamma_{12}^c=0.02\gamma_0$. Other parameters as in Fig. 2. For $r>100$ and $\Gamma_{12}^c=0$, the curves are almost identical to those of $r=100$ except for the actual pump intensities that decrease as r increases since $I_s\propto 1/r$

3. Reflectivity Behaviour

The main features of the resulting reflectivity behaviour are illustrated in Figs. 4 and 5 for several values of the attenuation-length parameter $\alpha_0 L$ after keeping the laser frequency fixed at $\Delta_0=0$ and choosing $r=100$, $\Gamma_{ij}^c=0$. Figure 4 shows the variation of the line center reflectivity with the pump intensity I_p and Fig. 5 the reflectivity spectra obtained by varying the sublevel splitting δ for two different equal-intensity pump irradiations. At low atomic densities and as typically observed in resonant DFWM, the line-center reflectivity peaks when the pump intensity is of the order of the saturation intensity, however the spectra of PC reflectivity using TOP as nonlinear generation mechanism are distinguished by single-peaked lineshapes since saturation affects equally the real and imaginary contributions of Zeeman coherence to the mixing susceptibility. Power broadening of the reflectivity resonance corresponding to optically thin media ($R \propto |K_2^{+}{}^*|^2 L$ for $\alpha_0 L \leqq 0.1$) is clearly seen in Fig. 5a and is basically described by the spectral halfwidth of the nonabsorption resonance (11).

The maximum obtainable reflectivity increases significantly when $\alpha_0 L$ is increased and in particular at pump irradiances in the saturation regime amplification is possible. Furthermore, for the case of exact resonance ($\delta=0$), the predominance of field-conjugate mixing over absorption as a consequence of coherent population trapping implies that PC reflectivity is dominated by a purely imaginary valued gain parameter. Reflectivity (4) behaves thus like a periodic function of $k/2$ GL attaining divergence when

$$\tan(GLk/2) = -G/X \tag{12}$$

if $(2p+1)\pi/2 < GLk/2 < (1+p)\pi$, as well as propagation induced annihilation when

$$GLk/2 = (1+p)\pi\,. \tag{13}$$

At a given large enough atomic density (G and X are proportional to N), the reflectivity versus pump intensity (Fig. 4) exhibits thus a series of alternating maxima and minima corresponding to the pump intensity values approaching the conditions (12, 13). Similar reflectivity behaviour may be observed in less favourable TOP conditions, as when r decreases or in case of finite values of Γ_{12}^c, although the pump intensity threshold for attainable oscillation increases and much higher atomic densities are needed since the increased saturation effects substantially reduce the mixing efficiency and thus the gain parameter (Fig. 3).

With an increase of the Zeeman splitting, laser detuning and resulting dispersive effects have an opposite influence on the weak counterpropagating beams and prevent coupled-mode oscillation. PC

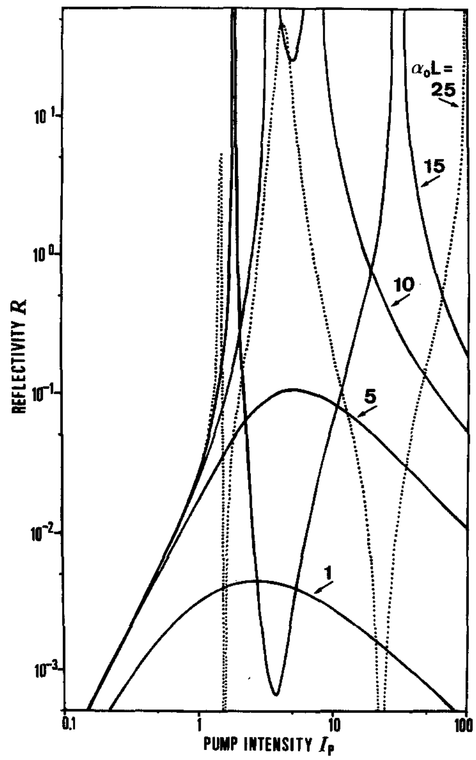

Fig. 4. Line center PC reflectivity as a function of the equal pump intensities $I_p = I_1 = I_2$, for several $\alpha_0 L$ values. Here $r=100$, $\Gamma_{ij}^c=0$, $ku=120\gamma_0$, and $\Delta_0=0$

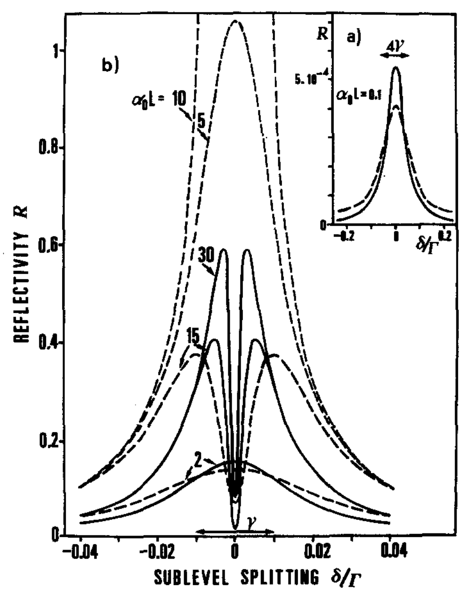

Fig. 5a, b. Phase-conjugate reflectivity spectra obtained by varying the normalized sublevel splitting $2\delta/\Gamma$ and keeping the laser frequency tuned to the degenerate one-photon-transition frequency i.e. $\Delta_0=0$. Two equal pump intensity irradiances are considered, one of $I_p=3$ (continuous line) and one of $I_p=6$ (dashed line). The remaining parameters are equal to those of Fig. 4. (a) Line shapes corresponding to an optically thin medium with $\alpha_0 L=0.1$. (b) Line shapes corresponding to optically denser media and for several increasing values of the attenuation-length parameter $\alpha_0 L$

amplification is consequently also restricted to a very narrow spectral region around $\delta=0$, upper limited by the power-broadened nonabsorption resonance half width. Of equivalent width is the propagation induced central dip appearing when at a given pump intensity $\alpha_0 L$ is further increased, after attainment of line-center oscillation, and approaches or fulfils the line-center reflectivity annihilation condition (13) (Fig. 5). The central dip is comparable to the typical saturation dip observed in previous DFWM experiments analysing TOP effects [5] in the sense that both occur in the narrow region where TOP is most effective and both reflect the dispersive contributions of the power-broadened Zeeman coherence resonance in the saturation regime. However, propagation effects play here an essential role and the dip is subjected to further narrowing as $\alpha_0 L$ or equivalently p in (13) increases.

The effect of nonabsorption resonance in the reflectivity behaviour is still very pronounced if the optical fields are out of resonance with the degenerate one-photon transition. The resulting line-center PC reflectivity as a function of $\alpha_0 L$ is shown in Fig. 6 for several laser frequencies tuned inside the one-photon Doppler profile ($0 \leqq \Delta_0 < ku$ and sublevel splitting parameter $\delta=0$). The total pump intensity is equal to that of Figs. 2 and 5, i.e. $I_1+I_2=6$, and Fig. 6a also corresponds to an equal-intensity pump irradiance while Fig. 6b illustrates the effect of pump imbalance for the particular case of a pump intensity ratio of I_1/I_2 or $I_2/I_1=2$. For Δ_0 of the order of the natural linewidth Γ ($\Delta_0 \leqq 2\Gamma$), noticeable Zeeman coherence is generated and since, in addition, the field absorption is effectively reduced as a consequence of the minor efficacy of the fields in exciting the optical transitions, PC oscillation results, without the requirements of an increase of the equal-intensity pump irradiance, at even lower atomic densities when compared to the exact resonant interaction. In contrast, the DFWM response arising from spatially dependent saturation of the population whether in two- or three-level systems gives rise to PC amplification or oscillation only when field detuning from one-photon resonance is large enough to reduce or even eliminate absorption but where, as a result from reduced interaction efficiencies, the generation mechanism begins only to show importance at very high pump intensities and atomic densities [10, 22]. In these laser detuned situations, differences in refractive index of the beams due to slightly unbalanced pump intensities, as, for instance, when one pump beam is obtained by retroreflection of the other, prevent oscillation. Here, when the laser frequency coincides with the one-photon transition frequency, pump asymmetry effects on the dispersion-free medium response (at $\delta=0$) do obviously not hinder the possibility of oscillation. Nevertheless, TOP is most effective when both saturating pump intensities are of equal magnitude and the diminution of the gain parameter as the pump intensity difference increases implies the requirement of an increase of the atomic density for attainment of oscillation. But when $\Delta_0 \neq 0$, dispersive effects on the weak beams due to laser detuning from the degenerate 0–1 and 0–2 transitions do not mutually cancel in the imbalanced pump irradiation case. As a result the maximum reflectivity decreases as the pump intensity ratio augments and, as typically observed in Doppler-broadened media, R has an anisotropic saturation behaviour such that PC generation is more favoured when the intensity of the backward pump I_1 is greater than forward pump intensity I_2.

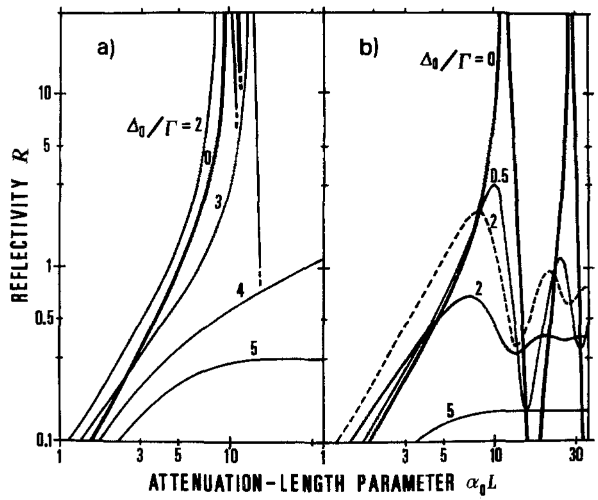

Fig. 6a, b. Phase-conjugate reflectivity as a function of the attenuation-length parameter $\alpha_0 L$ for several normalized values of the laser detuning Δ_0/Γ from the degenerate one-photon-transition frequency ($\delta=0$). The PC reflectivity corresponding to the particular case of $\Delta_0=0$ is represented by a thick line curve. (a) Equal pump irradiance of $I_1+I_2=6$. (b) Imbalanced pump irradiance of $I_1=2$, $I_2=4$ (thick and fine continuous line) and of $I_1=4$, $I_2=2$ (broken line). Other parameters as in Fig. 4

4. Conclusions

We have analysed the PC reflectivity behaviour arising from TOP in backward DFWM on inhomogeneously broadened coupled transitions. In the considered interaction pump waves as well as forward pump and probe wave have orthogonal polarizations and selectively interact with the different one-photon transitions. The main conclusion is that through the occurrence of nonabsorption resonance PC amplification and oscillation is possible on exact resonance with the resultant advantages of moderate optical densities, low pump intensities, much lower than those required to saturate the optical lines, and the admittance of pump imbalanced irradiation.

A more refined theory should take into account pump absorption and pump depletion effects [11]. For laser frequencies tuned inside the one-photon Doppler profile, pump absorption effects are actually not negligible. These effects are generally most important when the pump intensities are below the saturation intensity. However, high PC efficiencies are predicted in the saturation regime. Above the saturation intensity, pump irradiance within the four-wave mixer should be almost constant, and near to the saturation threshold, pump absorption effects are expected to shift somewhat the optimum pump intensity for PC amplification or oscillation. Pump depletion certainly plays an important role when PC reflectivity drastically increases and the generated wave intensity is a significant fraction of the pump intensity. Nevertheless the effect of depletion can be experimentally reduced by diminishing the incident probe intensity.

We wish to point out that the predicted reflectivity behaviour may be essentially met in typical experimental situations as for example in sodium vapors using cw pump irradiances of the order of milliwatts. Broadband laser irradiation should be also applicable for the obtention of very high PC reflectivities. Nevertheless, the considered DFWM interaction scheme does to our knowledge not appear to have been used previously to investigate experimentally the PC reflectivity.

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