Giant pulse lasing in three-level systems

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We propose an alternative method to Q switching for generating giant pulses of laser light in three-level media. This mehod is based on the presence of an external coherent field driving one transition to allow the accumulation of a large population inversion in the other transition without laser oscillation even in a cavity with high-Q factor. The switching off of the external coherent field causes the development of the giant pulse. Different time profiles for the switching off of the external field have been investigated.
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I. INTRODUCTION

Q switching is a standard technique for generating laser pulses of short duration, e.g., 10^{-7}–10^{-8} sec, and relatively high peak power, e.g., 10^6–10^7 W, the so-called giant pulses [1]. In this technique, the quality Q factor of the laser resonator is kept initially low to prevent laser oscillation while pumping builds up a large population inversion. Then a sudden switching of the cavity Q from the low value to a high value results in the release as a laser pulse of short duration of the energy accumulated in the upper level. The lifetime of this level must be relatively long, usually a fraction of a millisecond, to allow the accumulation of a large population inversion. This is the case of most crystalline solid-state lasers (e.g., Nd:YAG, ruby, alexandrite) and some gas lasers (CO_2 and iodine) [2]. In these systems the pumping mechanisms are incoherent and only *two levels* of the amplifying medium are coherently coupled with the light field.

We present here an alternative technique to Q switching for generating giant pulses of laser light in *three-level media*. In these configurations, the presence of an external coherent field acting on one transition modifies substantially the threshold population inversion needed for laser oscillation in the other transition. In fact, in the presence of an incoherent pumping mechanism in the latter transition, it is possible to accumulate a large population inversion without laser oscillation even in a lossless cavity. This phenomenon is known as population inversion without amplification or lasing (IWA or IWL) [3], one of the many known effects induced by atomic coherence [4]. After the inversion saturates to a high value, it is possible to extract the energy from the inverted transition in the form of a giant laser pulse by switching off the driving field in the other transition. Actually, in these three-level configurations the switching on/off of the driving field plays a similar role as the switching from a low to a high value of the cavity Q factor in two-level systems. We note that, to our knowledge, this is the first potential application of the inversion without amplification phenomenon.

The paper is organized as follows. In Sec. II the model for a laser operating on a cascade three-level system is introduced and explicit necessary conditions for the appearance of giant pulses are derived. In Sec. III we numerically investigate a specific example and compare the results to those of the conventional Q-switching technique. We study also the influence of different time profiles for the switching off of the external driving field. Some considerations for a practical realization of the new technique for generating giant laser pulses are presented in Sec. IV. Section V summarizes our main conclusions.

II. MODEL

We will focus our analysis to the cascade configuration shown in Fig. 1. A ring laser cavity is prepared in order to generate a laser field, the giant pulse, in the upper transition $|1\rangle$-$|2\rangle$ of the three-level media. The coupling of this field with the atomic transition is characterized by a Rabi frequency $2\alpha(t) \equiv \mu_{12}E_\alpha(t)/\hbar$ with μ_{12} the electric dipole moment of the $|1\rangle$-$|2\rangle$ transition and E_α the electric field am-

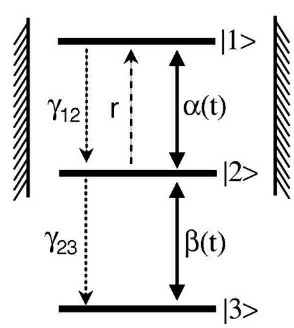

FIG. 1. Cascade three-level system under investigation. $\beta(t)$ and $\alpha(t)$ are half the Rabi frequencies of the external coherent and the generated pulsed laser fields, respectively. γ_{12} and γ_{23} are the spontaneous population decay rates of the laser and driven transitions and r is the population transfer rate of an incoherent pumping process in the laser transition.

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plitude of the laser pulse. In addition, an external pulsed laser field drives the lower transition $|2\rangle$-$|3\rangle$ with a Rabi frequency $2\beta(t)\equiv\mu_{23}E_{\beta}(t)/\hbar$. Without loss of generality, α and β are taken real. An incoherent continuous pumping process from level $|2\rangle$ to level $|1\rangle$, denoted by a rate r, models the mechanism needed to invert the population of the upper transition [5]. The population decay rates are given by γ_{12} and γ_{23} for the transition $|1\rangle$-$|2\rangle$ and $|2\rangle$-$|3\rangle$, respectively. The $|1\rangle$-$|3\rangle$ transition is allowed as a two-photon transition. In the standard density-matrix formalism with the rotating wave and slowly varying envelope approximations, the Maxwell-Schrödinger equations for our system read

$$\dot{\rho}_{11}=-\gamma_{12}\rho_{11}+r\rho_{22}+2\alpha y_{12}, \tag{1a}$$

$$\dot{\rho}_{22}=\gamma_{12}\rho_{11}-(\gamma_{23}+r)\rho_{22}-2\alpha y_{12}+2\beta y_{23}, \tag{1b}$$

$$\dot{\rho}_{33}=\gamma_{23}\rho_{22}-2\beta y_{23}, \tag{1c}$$

$$\dot{x}_{12}=-\Gamma_{12}x_{12}-\Delta_{\alpha}y_{12}+\beta y_{13}, \tag{1d}$$

$$\dot{y}_{12}=-\Gamma_{12}y_{12}+\Delta_{\alpha}x_{12}-\beta x_{13}+\alpha(\rho_{22}-\rho_{11}), \tag{1e}$$

$$\dot{x}_{23}=-\Gamma_{23}x_{23}-\Delta_{\beta}y_{23}-\alpha y_{13}, \tag{1f}$$

$$\dot{y}_{23}=-\Gamma_{23}y_{23}+\Delta_{\beta}x_{23}+\alpha x_{13}+\beta(\rho_{33}-\rho_{22}), \tag{1g}$$

$$\dot{x}_{13}=-\Gamma_{13}x_{13}-(\Delta_{\alpha}+\Delta_{\beta})y_{13}-\alpha y_{23}+\beta y_{12}, \tag{1h}$$

$$\dot{y}_{13}=-\Gamma_{13}y_{13}+(\Delta_{\alpha}+\Delta_{\beta})y_{13}+\alpha x_{23}-\beta x_{12}, \tag{1i}$$

$$\dot{\alpha}=-\kappa\alpha-gy_{12}, \tag{1j}$$

with $\rho_{ij}\equiv x_{ij}+iy_{ij}$ being x_{ij} and y_{ij} real variables. κ accounts for the cavity losses and $g\equiv\pi\nu_{12}N\mu_{12}^{2}/\hbar\varepsilon_{0}$ for the unsaturated gain parameter of the lasing transition with frequency ν_{12} and N being the density of atoms. Δ_{α} and Δ_{β} are, respectively, the laser field and driving field detunings from the corresponding atomic resonances. The laser field detuning is a variable linked with the cavity detuning Δ_{c} through the expression $\Delta_{\alpha}=\Delta_{c}-g(x_{12}/\alpha)$. Finally, the coherence decay rates are given by $\Gamma_{12}=(\gamma_{12}+\gamma_{23}+r)/2, \Gamma_{23}=(\gamma_{23}+r)/2$, and $\Gamma_{13}=\gamma_{12}/2$. In what follows, we will consider the completely resonant case, i.e., $\Delta_{c}=\Delta_{\alpha}=\Delta_{\beta}=0$, which means $x_{12}=x_{23}=y_{13}=0$ [6]. In our notation, $\alpha y_{12}>0(<0)$ and $\beta y_{23}>0(<0)$ mean, respectively, laser field and driving field absorption (amplification).

From a linear stability analysis of the trivial solution corresponding to $\alpha=0$, it is easy to show that the trivial solution can be destabilized leading to laser emission through either a pitchfork or a Hopf bifurcation [7]. The necessary and sufficient condition for the pitchfork bifurcation to occur is

$$n_{12}\equiv\rho_{11}-\rho_{22}>n_{\rm th}^{p}\equiv\frac{\kappa}{g}\left(\Gamma_{12}+\frac{\beta^{2}}{\Gamma_{13}}\right)+\frac{\beta y_{23}}{\Gamma_{13}}, \tag{2a}$$

while for the Hopf bifurcation it reads [8]

$$n_{12}>n_{\rm th}^{H}\equiv\frac{\Gamma_{12}+\Gamma_{13}}{g(\Gamma_{12}+\kappa)}[\kappa(\kappa+\Gamma_{12}+\Gamma_{13})+\Gamma_{12}\Gamma_{13}+\beta^{2}]-\frac{\beta y_{23}}{\Gamma_{12}+\kappa}. \tag{2b}$$

The smallest of these two population differences $n_{\rm th}^{p}$ and $n_{\rm th}^{H}$ determines the laser emission threshold. Clearly from Eqs. (2), there are two different ways of achieving a large inversion with no laser oscillation: (i) using a cavity with high losses κ (i.e., low-Q factor) and (ii) using a coherent driving field, which through the terms β^{2} and βy_{23} can greatly increase the lasing thresholds, giving rise to the so-called population inversion without lasing phenomena. In what follows we will focus on the second case in which inversion without lasing is due to the presence of the driving field. The two terms β^{2} and βy_{23} are related with the ac splitting of the common level $|2\rangle$ (Rabi sidebands or Autler-Townes doublet) and with the quantum interference between these sidebands. In our case in which there is population inversion, the gain spectrum consists of an Autler-Townes doublet with positive peaks at $\Delta_{\alpha}\approx\pm\beta$. However, it is well known [9,10] that the Autler-Townes doublet is not just the superposition of two Lorentzians centered at $\Delta_{\alpha}\approx\pm\beta$. These two resonances are not independent and therefore there is quantum interference between them. In our case this interference is destructive in the resonance region $-\beta<\Delta_{\alpha}<\beta$, which means that the gain there is smaller than that corresponding to the addition of two independent Lorentzians. As a result of this destructive quantum interference, one can even have negative gain (i.e., absorption) at $\Delta_{\alpha}=0$ in cases where there is population inversion. A detailed analysis of the influence of the quantum interference between the dressed-state resonances in different three-level atomic configurations can be found in Ref. [10]. Alternatively, inversion without amplification at line center in this configuration can be explained by using the quantum-jump formalism [9,10]. This formalism allows one to calculate the respective contributions of the various physical processes responsible for the amplification or the attenuation of the laser field: one-photon induced emission $|1\rangle\rightarrow|2\rangle$ and stimulated two-photon emission $|1\rangle\rightarrow|3\rangle$ or one-photon $|2\rangle\rightarrow|1\rangle$ and two-photon $|3\rangle\rightarrow|1\rangle$ absorptions, respectively. In our case one obtains [10] that at line center loss processes overcome gain processes, the main source of losses being two-photon absorption.

For the cascade scheme under consideration the driving field transition is not inverted, which means driving field absorption in the steady-state regime, i.e., $\beta y_{23}>0$. Thus, we will use the contribution of the term $\beta y_{23}/\Gamma_{13}$ to Eq. (2a) to prevent laser oscillation in the inverted transition through the pitchfork bifurcation even for a lossless cavity. On the other hand, from Eq. (2b) it is easy to see that even for the lossless cavity limit there is a threshold value for the gain parameter below which laser oscillation is not possible through the Hopf bifurcation. For $\kappa=0$ this threshold gain reads

$$g_{\rm th}(\kappa=0)=\frac{(\Gamma_{12}+\Gamma_{13})(\Gamma_{12}\Gamma_{13}+\beta^{2})}{\Gamma_{12}n_{12}+\beta y_{23}}. \tag{3}$$

Therefore, by choosing appropriate values for the parameters of the cascade system such that $n_{12}<\beta y_{23}/\Gamma_{13}$ [see Eq. (2a)]

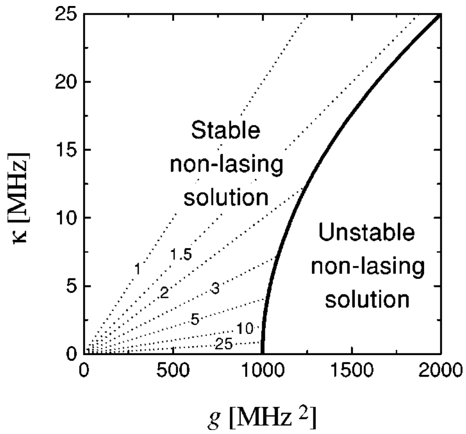

FIG. 2. Stability domain of the nonlasing solution in the κ-g plane for the following parameter setting: $\gamma_{12}=10$ kHz, $\gamma_{23}=50$ MHz, $r=50$ kHz, and $\beta=25$ MHz. At the left- (right-) hand side of the solid line, the nonlasing solution is stable (unstable). The dotted lines represent the relative amount of population inversion above the threshold population inversion for laser oscillation once the driving field is switched off.

and by taking the value of the gain parameter below $g_{\rm th}$, we can achieve a large inversion in the laser transition with no laser oscillation irrespective of the value of the cavity losses.

In a two-level system the nonlasing solution destabilizes through a pitchfork bifurcation and, by taking $\beta=0$ in Eq. (2a), the corresponding threshold population difference for laser oscillation reads $n_{\rm th}\equiv(\kappa/g)\Gamma_{12}$. In this case, a ''lossy'' cavity is used to achieve a large population inversion and the giant pulse is generated when the Q factor of the cavity is switched from a low value to a high value. In the three-level system under investigation, (i) we use the coherent field acting on the lower transition to achieve a large population inversion in the upper transition without laser oscillation, and (ii) we generate the giant pulse in the upper transition by switching off the driving field in the lower one. Once the driving field is switched off, the ratio $n_{12}/n_{\rm th}(\beta=0)$, being $n_{\rm th}(\beta=0)=(\kappa/g)\Gamma_{12}$, defines the relative amount of population inversion above the threshold population inversion for laser oscillation.

III. RESULTS AND DISCUSSION

Figure 2 shows the stability domain in the κ-g plane of the trivial nonlasing solution for the following parameter setting: $\gamma_{12}=10$ kHz, $\gamma_{23}=50$ MHz, $r=50$ kHz, and $\beta=25$ MHz. These relaxation rates are typical values for an electric-dipole forbidden, γ_{12}, and an allowed, γ_{23}, transition in the optical domain. Thus, as in traditional Q-switched systems, we assume an upper laser level lifetime in the range of a fraction of a millisecond [2]. The solid line represents the curve $n_{12}=n_{\rm th}^{H}$. At the left-hand side of the solid line $n_{12}<n_{\rm th}^{H}$, which guarantees the stability of the nonlasing solution since for these parameter values $n_{\rm th}^{p}\gg n_{\rm th}^{H}$, which means that the trivial solution can destabilize only through the Hopf bifurcation. For $g<g_{\rm th}(\kappa=0)=1000$ MHz^2, there is no laser oscillation irrespective of the value of the cavity losses. The dotted lines in Fig. 2 represent the relative amount of inversion above the threshold inversion just after

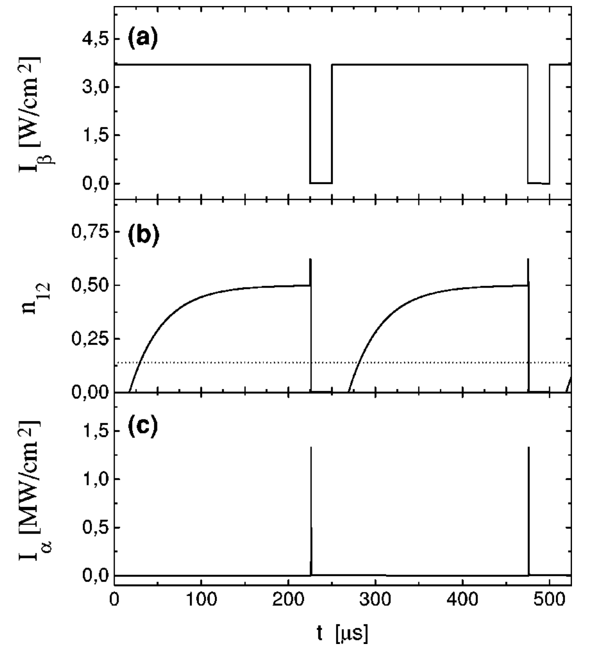

FIG. 3. Time evolution of the (a) driving field intensity, (b) population difference in the laser transition, and (c) intensity of the generated laser pulses. The parameters are as in Fig. 2 with $\kappa=5$ MHz, $g=900$ MHz^2, $\mu_{12}=10^{-31}$ C cm, and $\mu_{23}=10^{-28}$ C cm. The dotted line in (b) represents the threshold population inversion for laser oscillation in the absence of the driving field, i.e., $n_{\rm th}(\beta=0)$.

the switching off of the driving field, i.e., $n_{12}/n_{\rm th}(\beta=0)$. Figures 3 and 4 show the development of the giant laser pulse in the upper transition of the cascade scheme under investigation for the same parameters as in Fig. 2 with $\kappa=5$ MHz, $g=900$ MHz^2, $\mu_{12}=10^{-31}$ C cm, and $\mu_{23}=10^{-28}$ C cm. These values for μ_{12} and μ_{23} are consistent with the corresponding spontaneous population decays γ_{12}

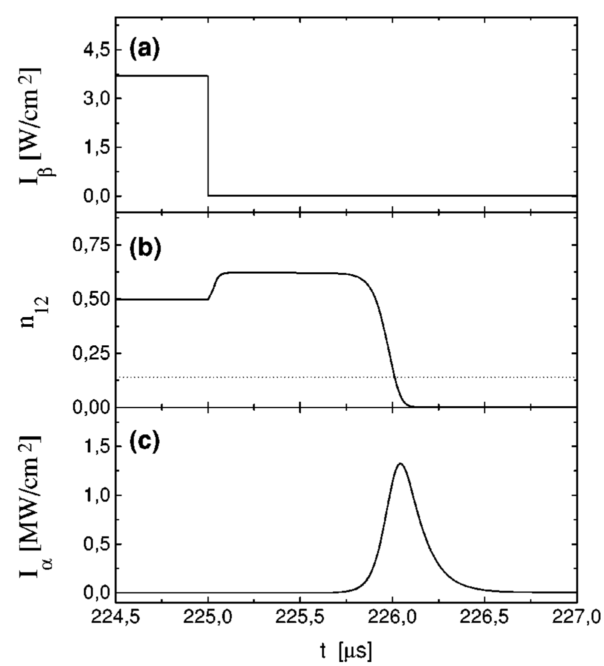

FIG. 4. Enlarged part of the time evolution of Fig. 3.

and γ_{23} provided that the transition frequencies satisfy $\omega_{12} \approx 6\omega_{23}$ [11]. Substituting these parameter values in Eqs. (1) with $\alpha=0$, we obtain the following steady-state population inversion with the driving field on $n_{12}^{ss}=0.5$. On the other hand, the threshold inversions for these parameters read from Eqs. (2a) and (2b): $n_{\mathrm{th}}^{p}=1193$ and $n_{\mathrm{th}}^{H}=0.61$. Thus, during the first 225 μs the presence of the driving field [Figs. 3(a) and 4(a)] with an intensity of 3.7 W cm^{-2}, which corresponds to $\beta=25$ MHz [12], prevents laser oscillation and a large population inversion accumulates [Figs. 3(b) and 4(b)] [13]. The dotted line in these figures represents the threshold population inversion in the absence of the external laser field, i.e., $n_{\mathrm{th}}(\beta=0)$. After the population inversion saturates, the driving field is switched off for 25 μs and a giant laser pulse develops since then the population inversion is well above the threshold population inversion (dotted line). For the parameters used, the peak intensity of the laser pulse gives $(I_\alpha)^{\mathrm{peak}}=1.33$ MW cm^{-2} with a pulse width at half maximum about 210 ns and an integrated pulse energy of 0.33 mJ cm^{-2}. Note that the pulse duration is similar to that of the longest pulses from traditional Q-switched systems [1]. The time delay between the instantaneous switching off and the pulse generation is 1.1 μs [14] and the repetition rate of the laser pulses is 4 kHz. It should be remarked that a laser field with a power of a few watts is able to control the generation of laser pulses of a few megawatts of peak power.

With the above parameters and $\omega_{12}=2\times10^{16}$ s^{-1}, $g=900$ MHz^2 requires an atomic density $N=8.5\times10^{13}$ cm^{-3}. Notice that for a given g the required density N is inversely proportional to $\omega_{12}\mu_{12}^2$. Therefore, by simply increasing γ_{12} (which scales as μ_{12}^2) by two orders of magnitude, the atomic density decreases by two orders of magnitude. Unfortunately, with $\gamma_{12}=1$ MHz and other parameters as for Fig. 3, the peak intensity of the laser pulses is quite small, $(I_\alpha)^{\mathrm{peak}}=12.3$ W cm^{-2}. One obtains the same atomic density $N=8.5\times10^{11}$ cm^{-3} by simultaneously decreasing the gain parameter by one order of magnitude to $g=90$ MHz^2 and increasing γ_{12} by one order of magnitude to $\gamma_{12}=0.1$ MHz. In this case, with $\kappa=0.5$ MHz, the peak intensity of the resulting pulses is $(I_\alpha)^{\mathrm{peak}}=2.5\times10^3$ W cm^{-2}. From the preceding results one can estimate that a lower bound for the atomic density required to observe the generation of laser pulses is $N\approx10^{12}$ cm^{-3}.

Figure 5 shows the intensity of the laser pulse, along pulse evolution, as a function of the relative amount of inversion above threshold for different values of the cavity losses κ. Clearly, decreasing the cavity losses, the peak intensity increases significantly since the amount of initial inversion also increases.

On the other hand, when the unsaturated gain parameter is increased from $g=900$ MHz^2 to $g=1100$ MHz^2, the nonlasing solution destabilizes through the Hopf bifurcation even in the presence of the driving field, since $n_{\mathrm{th}}^{H}(g=1100\ \mathrm{MHz}^2)=0.48<n_{12}^{\mathrm{ss}}=0.5$. In this particular case, after a transient the system reaches a cw regime with a laser field intensity $(I_\alpha)^{\mathrm{cw}}(g=1100\ \mathrm{MHz}^2)=2.7\times10^{-4}$ MW cm^{-2}. Notice that for this case, the ratio between $(I_\alpha)^{\mathrm{peak}}(g=900\ \mathrm{MHz}^2)$ and $(I_\alpha)^{\mathrm{cw}}(g=1100\ \mathrm{MHz}^2)$ is around 5000. In ordinary Q-switched systems the ratio between peak power of the pulse and steady-state power is also of this order of magnitude [15]. Finally, as a numerical example, for a resonator of length $l=15$ cm and a beam section of 0.1 cm^2, the giant pulse power outside the ring cavity can be easily calculated assuming a transmittance $T=2l\kappa/c=0.5\%$ to be $(P_\alpha)_{\mathrm{out}}^{\mathrm{peak}}=6.7$ kW. This is between two and three orders of magnitude smaller than the peak power that can be achieved with traditional Q-switched lasers [1,2]. However, it is perhaps more instructive to compare the above performance to that obtained in the case of applying the standard Q-switching technique to our system. For this purpose we substituted the coherent drive field by an incoherent continuous pump field characterized by a bidirectional transfer rate R between levels $|3\rangle$ and $|2\rangle$, and switched the cavity decay rate κ between $\kappa=1000$ MHz and $\kappa=5$ MHz. The cavity losses were set to the high κ value for time intervals of 225 μs, separated by time intervals of 25 μs with $\kappa=5$ MHz. For pump rates between $R=5$ MHz and $R=50$ MHz we always obtained peak powers smaller than in the previous case where coherent control of lasing was used [for instance, for $R=25$ MHz one obtains $(I_\alpha)^{\mathrm{peak}}=0.49$ MW cm^{-2}]. This is due to the absence of coherent coupling between pump and laser fields (one has now $\rho_{13}=0$) and also to the fact that while there was an increase in the population difference n_{12} after the coherent drive field was switched off [see Fig 4(b)], this increase obviously does not occur when the cavity Q is switched.

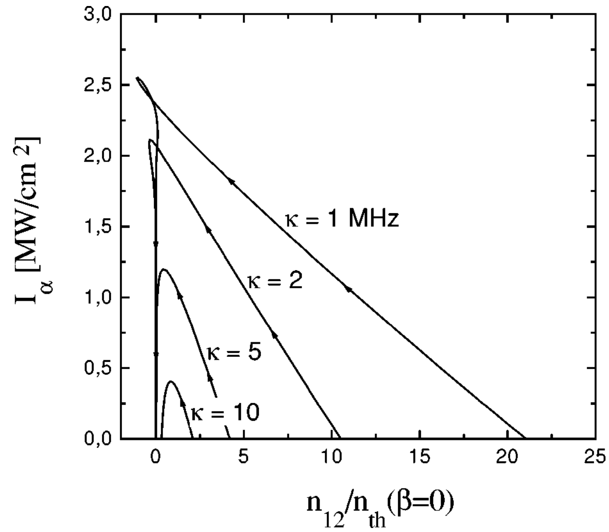

FIG. 5. Intensity of the generated laser pulses along time evolution as a function of the relative amount of population difference above the threshold population inversion for different values of the cavity losses κ.

In Fig. 6 the role of the time profile of the driving field is investigated by modeling its switching off through a hyper-Gaussian profile given by

$$I_\beta(t-t_0)=I_\beta^{(0)}e^{-c_n(t-t_0)^n} \tag{4}$$

with $c_n=3^{n-2}$ being $n=10,2,1.5,1.25,1,0.8$, the parameter that accounts for how fast the driving field is switched off, with $n=\infty$ corresponding to the square profile. $t_0=225$ μs indicates the switching-off time and $I_\beta^{(0)}=3.7$ W cm^{-2} the initial driving field intensity. As it is clearly seen in Fig. 6,

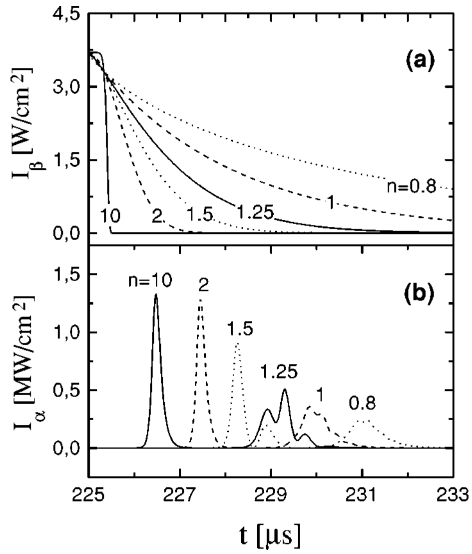

FIG. 6. (a) Different time profiles for the driving field switching off, and (b) the corresponding intensity of the generated laser pulses. Each curve is characterized by a different value of the parameter *n* that accounts for how fast the driving field is switched off [see Eq. (4) in the text].

the time delay until the pulse develops increases and the peak intensity decreases as we switch off slower the driving field intensity. For $n \leqslant 1.5$ the instantaneous population difference $n_{12}(t)$ crosses several times the instantaneous threshold population inversion $n_{\rm th}^H(t)$ and several pulses develop. For switching-off times larger than a few tens of μs the laser pulse almost vanishes.

It is worth remarking that up to now we have considered a homogeneously broadened system. The coherent control of lasing we have described is fragile with respect to Doppler broadening of the two-photon transition, in the most general case in which the frequencies of drive (ω_β) and laser fields (ω_α) are different. In fact, it is well known [16–18] that the modifications of the optical properties of the lasing transition, induced by the driving optical field, do not change significantly as long as the two-photon resonance condition $\Delta_\alpha = -\Delta_\beta$ is fulfilled. Thus, for $\Delta_\alpha = \Delta_\beta = 0$, we know that our system shows IWL in the absence of Doppler broadening. However, since the linear Doppler effect results in additional detunings, atoms with velocity $v \neq 0$ will ''see'' the frequencies of driving and laser fields ($\omega_\beta \neq \omega_\alpha$) affected by different Doppler shifts. These atoms will not fulfill the two-photon resonance condition and therefore will amplify the laser field, so that there will always be lasing.

As a guide for the practical development of the proposed method for generating giant pulses of laser light, let us summarize the main requirements: (a) cascade three-level configuration with population inversion at the lasing transition [19]; (b) metastable upper laser level ($\gamma_{12} \leqslant 10^4, 10^5$ Hz with $\mu_{12} \sim 10^{-31}$ C cm); (c) relatively high number densities $N \geqslant 10^{12}$ cm^{-3} (in the specific example investigated previously, $N \sim 10^{14}$ cm^{-3}); (d) Doppler-free configuration; (e) good laser cavity (κ less than a few MHz); and (f) fast driving field switching off (a switching time interval smaller than a few tens of μs). Conditions (e) and (f) are in principle easy to achieve.

IV. PRACTICAL CONSIDERATIONS

Although the study presented in this paper is intended to give a general insight into the phenomenon of coherent control of IWL and, eventually, of lasing, we give below some considerations regarding experimental systems in which the above ideas could be tested.

In our numerical example above we assumed $\beta/\gamma_{23} = 1/2$, which means that these driving field intensities can be easily attained with cw laser powers. Since the upper laser level is metastable (as in conventional *Q*-switched systems), the needed incoherent excitation rates are also very moderate, $r > \gamma_{12} \leqslant \gamma_{23}$.

The impurity ions in media used in crystalline solid-state lasers (e.g., ruby or Nd:YAG) have homogeneously broadened electric dipole forbidden transitions between states arising from the inner unfilled shells [2]. Conditions (b), (c), and (d) above are fulfilled and it would be necessary to explore the energy level diagrams of these media to know whether there are suitable cascade configurations to fulfill condition (a).

A practical candidate to test the proposed method is CO_2. Population inversion in the regular bands $00^01 - 10^00(02^00)$ of CO_2, centered at 10.4 μm (9.4 μm), can be efficiently achieved by electrical pumping in low pressure CO_2:N_2:He mixtures. The lower level $10^00(02^00)$ of the lasing transition can be coupled by means of a HCl laser at 14 μm (by means of an optically pumped CF_4 laser at 16 μm) [20] to the 01^10 level for the coherent control of the laser emission at 10.4 μm (9.4 μm). Due to the long wavelengths involved, the Doppler broadening plays a minor role and the decay of the various levels is essentially determined by collisions, leading to lifetimes of the order of tenths of a millisecond. Thus, this system fulfills conditions (a), (b), (c), and (d) above.

Rydberg levels of alkali-metal atoms with principal quantum numbers $n > 20$ are long lived since the lifetime scales approximately as n^3. Rubidium is particularly interesting since many highly excited *S* and *D* states can be reached from the ground *S* state with four diode lasers. The main problem with Rb is Doppler broadening. Possible remedies are to make the driving Rabi frequency comparable to the Doppler width [16–18] or to add a buffer gas to use the Dicke narrowing of the two-photon transition for the suppression of its Doppler broadening. Unfortunately, in addition to the high power needed in both cases, it decreases the gain itself too, which might be compensated by the increase of the density of atoms in a cell. An alternative approach could be to use a dense atomic beam, but it is important to notice that the required atomic densities are difficult to achieve in practice. Although technically more difficult, the best solution is to use a sample of Rb atoms that has been laser cooled (thus eliminating Doppler broadening) and then eventually subjected to magnetic trapping and evaporative cooling to increase the atomic density.

V. CONCLUSIONS

We have proposed a method to generate giant pulses of laser light in three-level media based on the inversion without amplification phenomenon. In this technique, an external coherent field couples one transition while an incoherent pumping mechanism acting on the other transition creates a large population inversion. The presence of the external laser field prevents laser oscillation in the inverted transition even in a high-Q cavity on resonance with this transition. The switching off of the external coherent field results in the release as a laser pulse of short duration of the energy accumulated in the laser transition. In particular, we have considered a cascade three-level system deriving explicit necessary conditions for the appearance of the giant laser pulses. We have shown that an external coherent field with a power of a few watts is able to control the generation of laser pulses of a few megawatts of peak power. Different time profiles for the switching off of the external coherent field have been investigated showing that several laser pulses can develop when the driving field is adiabatically switched off. A comparison between the expected performances of the proposed scheme and those of traditional Q-switched systems has been performed and some considerations for a practical realization of the technique have been discussed.

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[4] For reviews on the topic see M. O. Scully, Phys. Rep. **219**, 191 (1992); E. Arimondo, in *Progress in Optics XXXV*, edited by E. Wolf (Elsevier Science, Amsterdam, 1996); S.E. Harris, Phys. Today **50**(7), 36 (1997), and references therein.

[5] Note that to invert the population between levels $|1\rangle$ and $|2\rangle$, the incoherent pumping scheme must consist of a transfer of population from level $|2\rangle$ to an atomic level above level $|1\rangle$ (not shown in Fig. 1), followed by a radiative or nonradiative decay to level $|1\rangle$. As in most theoretical treatments, this decay is considered here so fast that the population of the extra level can be neglected.

[6] In the numerical simulations we will maintain small but nonzero detunings to allow dispersive effects to take place and, therefore, including the possibility of detuned continuous wave lasing even with driving and cavity on resonance with their respective transitions. See R. Corbalán, J. Mompart, R. Vilaseca, and E. Arimondo, Quantum Semiclassic. Opt. **10**, 309 (1998); and Refs. [7] and [8].

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[8] Depending on the parameter values, the Hopf bifurcation yields to a continuous wave (cw) or a self-pulsing lasing regime, see Ref. [7] and A.G. Vladimorov, P. Mandel, S.F. Yelin, M.D. Lukin, and M.O. Scully, Phys. Rev. E **57**, 1499 (1998).

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[11] From the expression of the Einstein A coefficient, one obtains the following relationship between the spontaneous decay rates: $(\gamma_{12}/\gamma_{23})=(\mu_{12}/\mu_{23})^2(\omega_{12}/\omega_{23})^3$.

[12] The relationship between the laser intensity I and the Rabi frequency Ω is given by $I=\frac{1}{2}\epsilon_0 c\hbar^2\,(\Omega/\mu)^2$. When I is expressed in W cm^{-2}, Ω in Hz, and μ in C cm, $I=1.48\times10^{-71}\,(\Omega/\mu)^2$.

[13] In the numerical simulation, we fix a minimum Rabi frequency $\alpha^{(0)}=10^{-3}$ MHz which corresponds to an intensity $I_\alpha^{(0)}=5.92\times10^{-3}$ W cm^{-2} in order to allow, if possible, the destabilization of the trivial nonlasing solution.

[14] This time delay depends on the minimum Rabi frequency $\alpha^{(0)}$. For values of $\alpha^{(0)}/\gamma_{13}$ in the range $(10^{-4},10)$ this dependence reads as follows: Time delay$=T_0+k\ln(\alpha^0/\gamma_{13})$, with $T_0=0.815$ μs and $k=-0.097$ μs. This means that a large change in the value of $\alpha^{(0)}$ is practically equivalent to a shift in time scale by a small amount. Note that with $\alpha^{(0)}=0$ the time delay would be infinite, i.e., the laser does not switch on. This is a feature of the semiclassical laser equations, which ignore spontaneous emission into the laser mode [see Eq. (1j)].

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[19] For an on-resonance driving field, inversion without amplification can also be obtained in the so-called V-type and Λ-type three-level systems provided that there is also inversion in the driven transition (see Ref. [10]). As in the cascade configurations, absorption is due to the predominance of the two-photon loss processes at line center.

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