Coherent population transfer in Rb atoms by frequency-chirped laser pulses

G. P. Djotyan, J. S. Bakos, G. Demeter, P. N. Ignácz, M. Á. Kedves, Zs. Sörlei, J. Szigeti, and Z. L. Tóth
KFKI-Research Institute for Particle and Nuclear Physics, P.O.B. 49, Budapest H-1525, Hungary
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We investigate the behavior of ^{85}Rb atoms in the field of a sequence of frequency-chirped short laser pulses. The analysis is based on a numerical solution of equations for the probability amplitudes of the hyperfine levels of the $5S_{1/2}$-$5P_{3/2}$ transition in the ^{85}Rb atom and the dressed-states analysis. We analyze different regimes of interaction, including relatively short laser pulses (when the width of the pulse envelope spectrum is of the order of or exceeds the frequency interval between the hyperfine levels resulting in effective mixing of them) and relatively long ones (when the ground hyperfine levels are resolved but the excited ones are not resolved). In the latter case dependence of the population transfer efficiency on the initial coherence of the ground states is analyzed. The case of long laser pulses when all working hyperfine levels are resolved is also discussed using numerical simulations and a dressed-states analysis. We show that in all regimes considered, the interaction of a frequency-chirped laser pulse with the multilevel ^{85}Rb system is similar to the interaction with an effective two-level atom at sufficiently large peak intensities of the pulses. It allows us to perform efficient excitation of the multilevel atom by transferring populations of two hyperfine ground states to the excited ones and back to the ground states using a pair of frequency-chirped laser pulses. We propose to utilize this scheme of population transfer for the coherent manipulation of a beam of ^{85}Rb atoms using sequences of counterpropagating frequency-chirped short laser pulses.

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I. INTRODUCTION

Coherent manipulation of atoms without heating them in a process of changing their velocity or position is an important problem to be addressed in atomic interferometers, during the transportation of atoms between different stages of a laser-cooling system, and in other applications where conservation of the atomic phase is an important issue [1–3]. The deflection and splitting of atomic beams using trapped (dark) atomic states was realized in [4]. Coherent manipulation of atomic beams was proposed and demonstrated in counterpropagating laser beams in the scheme of stimulated Raman adiabatic passage (STIRAP) (see [5] and references therein) and also in multilevel Zeeman systems [6]. An important advantage of the STIRAP scheme is that the population transfer between ground states takes place without considerable population of the intermediate excited state, which preserves coherence of the interaction. The atomic recoil technique in standing wave geometry also presents a possibility for manipulation of atomic beams [7] (see also [8] for bichromatic counterpropagating laser beams), although a large frequency detuning is necessary to avoid spontaneous emission for preserving the atomic coherence.

Another way of coherent manipulation of atomic beams is stimulated excitation and de-excitation of atoms by counterpropagating pairs of short frequency-chirped laser pulses in the adiabatic passage (AP) regime of interaction [9–12]. It may be easily understood in the case of a two-level atom. First, a frequency-chirped laser pulse from one direction excites the atom and transfers mechanical momentum equal to $\hbar k_p$ (with k_p being the wave number of the laser pulse) in the direction of the laser pulse propagation. Another frequency-chirped laser pulse applied to the atom from the opposite direction with a time delay much shorter than the decay time of the excited state moves the atomic population to the ground state and transfers another $\hbar k_p$ momentum in the direction opposite to the laser pulse propagation direction. As a result the atom receives mechanical momentum equal to $2\hbar k_p$ after the action of the pair of the counterpropagating frequency-chirped laser pulses. The conditions for the realization of the AP regime in a two-level system are discussed in Ref. [13] (for more complicated multilevel systems see Ref. [5]). Duration of the pulses along with the time interval between them have to be much shorter than the relaxation times of the atomic system to avoid stochastization (and lost of the coherence) due to the spontaneous decay. Note that the stimulated force on the atoms applied by the counterpropagating laser pulses may be by orders of magnitude stronger than the one that resulted in the Doppler scheme of manipulation of atomic beams [10–12]. Another advantage of this scheme (compared with, for example, the STIRAP scheme) is the large number of the laser pulses (and, correspondingly, the larger mechanical momentum transferred to atoms) interacting with the atomic beam. This number of laser pulses is defined by the repetition rate of the laser source and the time of flight of the atoms through the interaction region.

In this paper we analyze the coherent interaction of frequency-chirped laser pulses with ^{85}Rb atoms. The Rb atom is an often-used model system in investigations in the field of laser cooling and manipulation of neutral multilevel atoms. While coherent manipulation of two-level atoms by counterpropagating laser pulses is well understood, there are much less results on coherent laser manipulation of multilevel atoms. Note here the only two examples where coherent manipulation of multilevel atoms is required: the atomic interferometers using multilevel atoms, and the problem of moving the multilevel atoms confined in the magneto-optic trap to the next step of the laser cooling in the technique of the Bose-Einstein condensate production. As we mentioned above, an effective coherent manipulation of the atoms minimizing the heating effect of the spontaneous transitions may be achieved by acting on the atomic system by pairs of counterpropagating frequency-chirped laser pulses. The aim of this investigation is to develop a simple, effective, and robust scheme for the coherent transfer of the population in a Rb atomic model from the hyperfine electronic ground states to the manifold of excited hyperfine states and back to the ground states by acting with pairs of frequency-chirped laser pulses. This process of the population transfer differs from those in the STIRAP schemes, which utilize trapped atomic states and delayed, partially overlapping laser pulses in the "counterintuitive" sequence to transfer atomic population between ground states without considerable excitation of the atom. On the contrary, we have the excitation of the Rb atom. To avoid the heating of the atomic beam due to the spontaneous decay of excited states, we assume that the duration of the laser pulses along with a time interval between (not overlapping) pulses forming the pulse pairs are much shorter than the relaxation times of the atomic system. Note that we need no initial optical pumping of atomic populations into one of the ground states as in the STIRAP scheme: as our analysis shows (see below), effective excitation and deexcitation of the atom by frequency-chirped laser pulses is mostly insensitive to the initial distribution of the atomic population between the ground states.

We investigate different regimes of interaction including relatively short and relatively long laser pulses. In the case of short laser pulses, the width of the pulse envelope spectrum is of the order of or exceeds the frequency interval between the hyperfine levels that results in the effective mixing of them. We refer to this case as "broadband pumping." On the contrary, the two ground hyperfine levels of the Rb model atom are resolved in the case of relatively long pulses with a relatively narrow width of the pulse envelope spectrum referred to as "narrow band pumping."

The working level system of the ^{85}Rb atoms is chosen to be its hyperfine levels manifold of the $5S_{1/2}$-$5P_{3/2}$ transition. The decay time of the excited level $5P_{3/2}$ is $\tau = 27$ ns [14]. The interaction with the atom to be coherent, the pulse duration must be shorter than relaxation times of the atomic system. The laser pulses with duration $\tau_p \leqslant \tau_c$, where the limiting duration τ_c is of the order of 2–3 ns, seem to be appropriate for the interaction to be coherent in the case under consideration. One can estimate that the corresponding width Δf_p of the laser pulse intensity envelope spectrum ($\Delta f_p \propto 1/\tau_p$) is smaller than the frequency interval ω_{12} ($\omega_{12} = 3035.7$ MHz, see Fig. 1) between the two ground-state hyperfine levels of the ^{85}Rb atom for laser pulse duration in the region of $1/\omega_{12} < \tau_p < \tau_c$. This case is referred to as narrow band pumping in this paper. For pulse duration $\tau_p < 1/\omega_{12}$, the laser pulse intensity envelope spectrum exceeds the spectral distance between the two ground states: $\Delta f_p > \omega_{12}$. It corresponds to the case of broadband pumping. Note that Δf_p exceeds the frequency distance between the excited hyperfine levels (see Fig. 1) for both narrow band and broadband pumping when $\tau_p < \tau_c$. Because the excited hyperfine levels manifold is not resolved in this case, the conditions of the adiabaticity of transitions to individual hyperfine levels

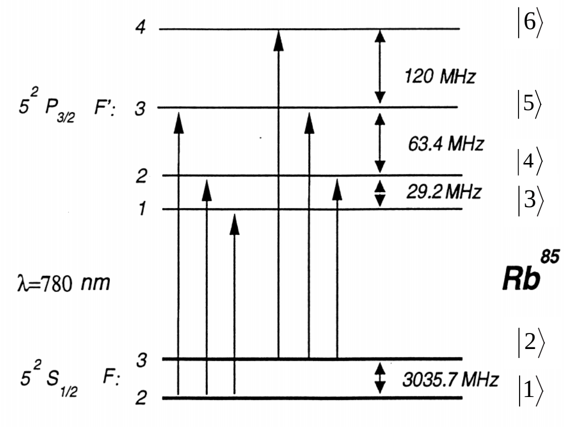


FIG. 1. The scheme of the working hyperfine levels of the ^{85}Rb atom.

cannot be fulfilled. To analyze these two cases of the broadband and narrow band pumping, we performed a numerical solution of the Schrödinger equations for the atomic probability amplitudes.

Quantum interference effects are expected to contribute to the population transfer process from the two not resolved ground hyperfine states in the case of broadband pumping when initially, the atomic population is present in both ground states. The analysis below shows that excitation efficiency of individual excited hyperfine levels depends on the phase difference of the ground state probability amplitudes (the initial coherence between the ground states). A similar situation was investigated in Ref. [15] for Raman amplification in a multilevel system.

To have a complete picture for the interaction of a frequency-chirped laser pulse with a multilevel system, we simulated also a case of long laser pulses when the hyperfine levels of the Rb model atom are resolved for both the ground and excited manifolds of the hyperfine levels assuming that conditions for the AP regime of interaction has been fulfilled. The results of the numerical solution of the Schrödinger equation are in good agreement with the results of the dressed-states analysis. While the conditions of the AP regime of interaction have a relatively simple form in the case of a two-level atom in the field of a frequency-chirped laser pulse [12], the situation is more complicated in the case of the multilevel quantum system under consideration. A laser pulse with a chirped frequency encounters a combination of different configurations (Λ,V) of quantum subsystems during interaction with the manifold of hyperfine levels of the ^{85}Rb atom's $5S_{1/2}$-$5P_{3/2}$ transition. A case when a chirped laser pulse couples a single ground state with a manifold of closely located excited states was discussed in Ref. [16]. It was shown that depending on the direction of the frequency chirp, the population is transferred into the lowest or the highest excited state. Note that even in the case of the AP regime of interaction, the quantum system under consideration in this paper is more complex than the one considered in Ref. [16] because of the presence of two ground states.

One of the main results of the analysis performed in this paper is that the interaction of a sequence of frequency chirped laser pulses with the multilevel ^{85}Rb system in all cases mentioned above leads to the effective excitation and deexcitation of the atom at sufficiently high peak intensities of the frequency-chirped laser pulses.

The remainder of this paper is organized as follows. In Sec. II we present the mathematical formalism for the analysis of a frequency-chirped laser pulse interaction with a six-level quantum system modeling the $5S_{1/2}$-$5P_{3/2}$ transition in the ^{85}Rb atom. The results of numerical simulations are presented in Sec. III for the broadband and narrow band laser pulses. In Sec. IV we present the results of numerical simulations and the dressed-states analysis for the case of a frequency-chirped laser pulse of sufficiently long duration when all the working hyperfine levels are resolved, assuming that conditions for the AP regime of interaction are fulfilled. A discussion and summary of the main results obtained in this paper are presented in Sec. V.

II. THE MATHEMATICAL FORMALISM

We begin with the Schrödinger equation for the probability amplitudes a_j, $j=1,\ldots,6$, of the atomic states of the ^{85}Rb atom (see Fig. 1) interacting with a frequency-chirped laser pulse having duration much shorter than all relaxation times of the atomic system. The envelope of the laser pulse amplitude is taken as a Gaussian in simulations below: $A(t)=A_0\exp[-t^2/2\tau_L^2]$ with $2\tau_L$ being the full duration (at the e^{-1} level) of the laser pulse intensity $I(t)\propto|A(t)|^2$. A linear chirp of the carrier frequency is assumed for simplicity: $\omega_L=\omega_{L0}+\beta t$, with ω_{L0} and β being the central frequency and the chirp velocity.

The Schrödinger equation for the state vector $\underset{\sim}{a}=(a_1,a_2,a_3,a_4,a_5,a_6)^T$ has the form

$$\frac{d}{dt}\underset{\sim}{a}=i\hat{H}\underset{\sim}{a}, \tag{1}$$

and the rotating-wave-approximation Hamiltonian is

$$\hat{H}=\begin{pmatrix} -\varepsilon_1 & 0 & \Omega_{13} & \Omega_{14} & \Omega_{15} & 0 \\ 0 & -\varepsilon_2 & 0 & \Omega_{24} & \Omega_{25} & \Omega_{26} \\ \Omega_{13}^* & 0 & 0 & 0 & 0 & 0 \\ \Omega_{14}^* & \Omega_{24}^* & 0 & -\varepsilon_4 & 0 & 0 \\ \Omega_{15}^* & \Omega_{25}^* & 0 & 0 & -\varepsilon_5 & 0 \\ 0 & \Omega_{26}^* & 0 & 0 & 0 & -\varepsilon_6 \end{pmatrix}, \tag{2}$$

where $2\Omega_{ij}=A(t)d_{ij}/\hbar$ is the time-dependent Rabi frequency and d_{ij} is the dipole moment for transition between states $|i\rangle$ and $|j\rangle$. $\varepsilon_1(t)=\Delta_0+2\beta t-\omega_{12}$ is the detuning of the time-dependent laser frequency from the one-photon resonant transition frequency between the states $|1\rangle\leftrightarrow|3\rangle$ (see Fig. 1) with ω_{12} being the frequency distance between the ground states $|1\rangle$ and $|2\rangle$. Δ_0 is the detuning of the central (at $t=0$) laser frequency from the resonant frequency for the (forbidden) transition between the states $|2\rangle$ and $|3\rangle$. The time-dependent detuning $\varepsilon_2(t)=\Delta_0+2\beta t$. Δ_0 is assumed to be zero in our analysis for simplicity. The constant frequencies $\varepsilon_4,\varepsilon_5,\varepsilon_6$ are equal to the frequency distance of the states $|4\rangle,|5\rangle,|6\rangle$ from the state $|3\rangle$, respectively (see Fig. 1).

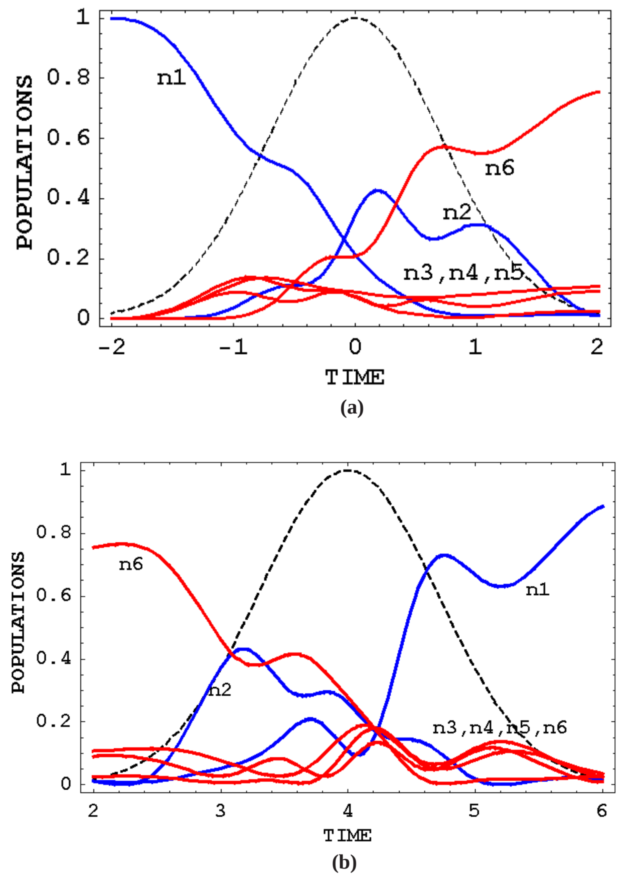


FIG. 2. (Color online) The population dynamics for the case of broadband pumping. Time is normalized by the duration of the laser pulse. The atom is assumed in a single ground state initially. The dashed line is the normalized intensity profile of the laser pulse. Time is normalized by the duration of the laser pulse. The dashed line is the normalized intensity profile of the laser pulse. The parameters applied are as follows: the duration of the pulse intensity $2\tau_p=10^{-1}$ ns; the peak Rabi frequency has values between $\Omega_R=42$ GHz and $\Omega_R=120$ GHz for allowed transitions between the two ground and four excited states; the chirp speed $\beta=200$ GHz/ns. (a) During the first pulse. (b) During the second pulse.

III. RESULTS OF THE NUMERICAL SIMULATIONS

A. Broadband pumping

First, we analyze the case of the broadband pumping when the width $\Delta f_p\propto 1/\tau_p$ of the envelope spectrum of the laser pulses exceeds the frequency distance ω_{12} between the two ground states and distances $\omega_{34},\omega_{45},\omega_{56}$ between the excited hyperfine levels of the atom (see Fig. 1): $\Delta f_p>\omega_{12},\omega_{34},\omega_{45},\omega_{56}$. The hyperfine levels population dynamics in the field of the frequency-chirped laser pulse is shown in Fig. 2 as the result of the numerical solution of Eq. (1). We assume that the population of the atom initially (without laser pulse) is concentrated in a single ground state.

As it can be seen from Fig. 2(a), the population of the ground states is completely transferred to the manifold of the excited states as a result of the interaction with the first chirped laser pulse. We use the values of the populations of the atomic states at the end of the laser pulse as initial conditions for the interaction with the second laser pulse having the same parameters as the first one. The dynamics of the quantum state populations during the action of the second laser pulse is shown in Fig. 2(b). As it can be seen, all population of the atom is again collected in the ground states at the end of the second laser pulse.

Two excited hyperfine levels $|4\rangle$ and $|5\rangle$ of the ^{85}Rb atom are commonly connected to both ground states $|1\rangle$ and $|2\rangle$ (see Fig. 1). One can expect quantum interference effects in transitions from the two ground states to the excited $|4\rangle$ and $|5\rangle$ states through different quantum passes in the field of the broadband laser pulse when the atomic population is located in both ground states initially. The population transfer in this case has to depend on the initial relative phase of the probability amplitudes of the ground states and on the initial population distribution among these states. To stress and reveal these interference effects, we use a ^{85}Rb atomlike model in which the same values for transition dipole moments are assumed for transitions between the ground and excited hyperfine levels. Also, we assume the same initial population of the ground states. As it follows from Fig. 3(a), the levels $|4\rangle$ and $|5\rangle$ are not populated at all after interaction with the frequency-chirped broadband laser pulse in the case when the ground states $|1\rangle$ and $|2\rangle$ are equally populated initially but have equal to π radian phase difference $\Delta\phi_0$ between their initial probability amplitudes. This is a result of destructive interference between common quantum transitions from the two ground states induced by the broadband pumping. It is important to note that even in this case we have complete transfer of the ground-state population to the excited states $|3\rangle$ and $|6\rangle$ having no common transitions to the two ground states. On the contrary, when the phase difference $\Delta\phi_0$ is assumed to be zero, all excited states are populated (with predominantly population of the states $|4\rangle$ and $|5\rangle$ due to the constructive interference) as a result of interaction with the laser pulse in Fig. 3(b). Note that all atomic population is transferred from the ground states to the excited states for both limiting values of the phase $\Delta\phi_0$. Numerical analysis shows that similarly, all population of the atom is being transferred to the ground states as a result of interaction with a second frequency-chirped laser pulse.

B. Narrow band pumping

In the case of narrow band pumping, $\omega_{34},\omega_{45},\omega_{56}<\Delta f_p<\omega_{12}$, only the ground state levels are assumed to be resolved when the levels of the excited manifold are mixed because of the relatively wide laser pulse envelope spectrum. The dynamics of the population transfer for the case of narrow band pumping is shown in Fig. 4. As can be seen, a complete population transfer between the ground and excited states may be achieved in this case too: the atomic population is completely transferred to the excited states after the first frequency-chirped pulse and is transferred back to the ground states as a result of interaction with the second laser pulse.

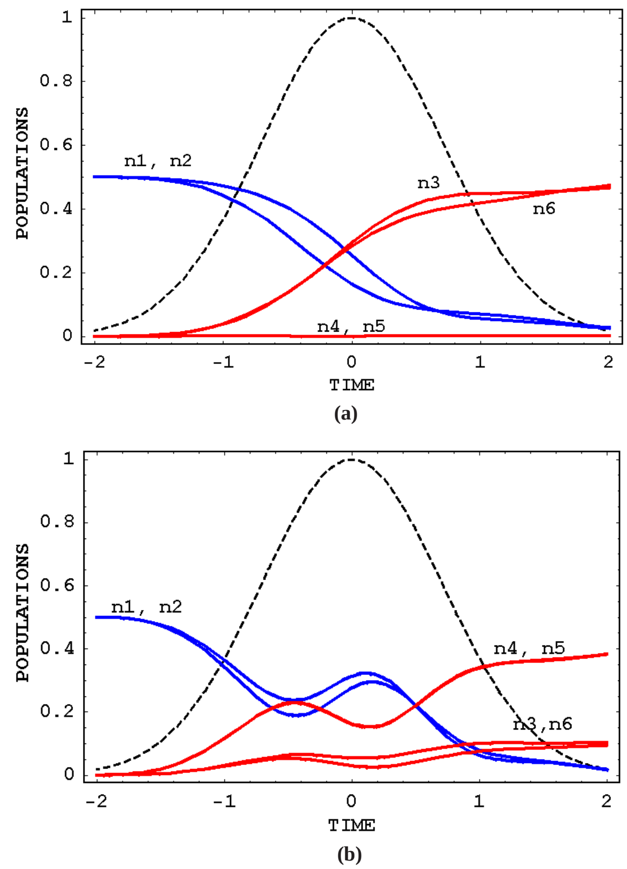
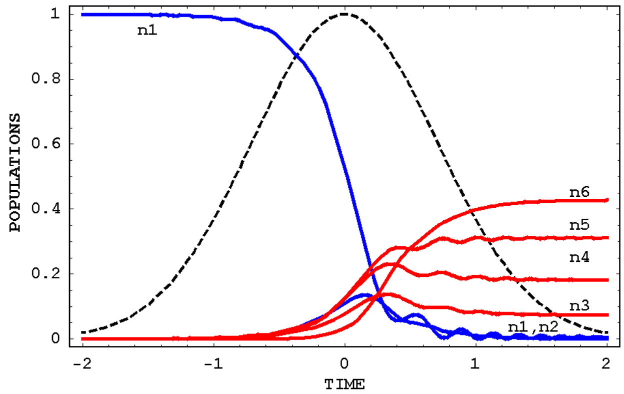


FIG. 3. (Color online) The population dynamics for broadband pumping. The two ground levels of the atom are equally populated initially. Time is normalized by the duration of the laser pulse. The parameters applied are as follows: the duration of the pulse (intensity) $2\tau_p=10^{-1}$ ns; the peak Rabi frequency is assumed to be the same and equal to $\Omega_R=40$ GHz for all allowed transitions between the two ground and four excited states; the chirp speed $b=140$ GHz/ns. The initial phase difference of the ground states probability amplitudes is (a) $\Delta\phi_0=\pi$ and (b) $\Delta\phi_0=0$.

The peak Rabi frequencies of transitions from the ground states to the excited ones used in the simulations exceeds the interval between the ground states. To examine the influence of power broadening of the hyperfine levels on the efficiency of the population transfer, we increased artificially the interval ω_{12} between the two ground states in order for ω_{12} to exceed significantly the peak Rabi frequencies of the transitions. The efficiency of the population transfer did not change significantly compared to the previous case of the original interval between the ground-state levels. It means that even in the case of the Rabi frequency exceeding the frequency interval between the hyperfine ground (or excited) levels, they could not be interpreted as mixed states in the interaction process with frequency-chirped laser pulses. Note that narrow-band excitation and spectral selectivity preservation were achieved with broadband lasers with chirped frequency even when the Rabi frequency exceeded the distance



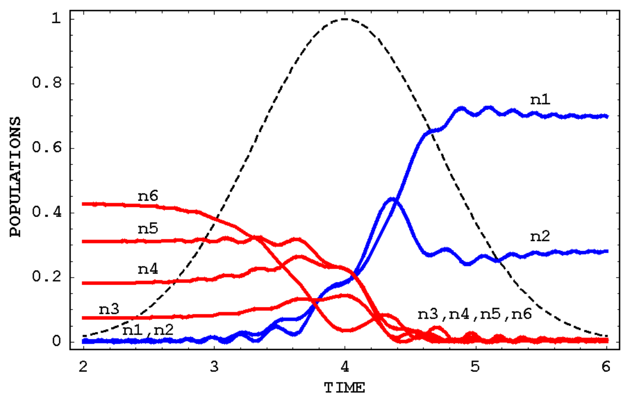


FIG. 4. (Color online) The population dynamics for the case of narrow band pumping. Time is normalized by the duration of the laser pulse. The atom is assumed in a single ground state initially. The dashed line is the normalized intensity profile of the laser pulse. The parameters applied are the duration of the pulse (intensity) $2\tau_p=1$ ns; the peak Rabi frequency is between $\Omega_R=3.15$ GHz and $\Omega_R=9$ GHz for the allowed transitions between the two ground and four excited states; the chirp speed $b=7$ GHz/ns. (a) During the first pulse. (b) During the second pulse.

between the atomic levels [16] or the width of the Doppler-broadened lines of the atomic ensemble [17].

To make the picture of the interaction of ^{85}Rb model atoms with frequency-chirped laser pulses even more complete, we proceed to examine this interaction for a case when the hyperfine levels are resolved for both the ground states and the excited state manifold using the dressed-states analysis in the next section. The results obtained in the dressed-states picture are compared with the results of the numerical solution of the Schrödinger equation for the same quantum system.

IV. THE DRESSED-STATE PICTURE

The dressed-state analysis [18,19] is a convenient way to predict the behavior of the atomic populations as a result of interaction with the frequency-chirped laser in the AP regime. The solution $\underline{a}(t)$ of Eq. (1) can be represented in the basis of the adiabatic dressed states $\underline{b}^{(k)}(t)$,

$$\underline{a}(t)=\sum_k r_k(t)\underline{b}^{(k)}(t)\exp\left[-i\int_{-\infty}^{t}w_k(t')dt'\right],$$

with the initial condition at $t\rightarrow-\infty$,

$$\underline{a}(-\infty)=\sum_k r_k(-\infty)\underline{b}^{(k)}(-\infty),$$

where $\underline{b}^{(k)}(t)$ is the dressed-state eigenvector, corresponding to the kth eigenvalue (quasienergy) w_k of the Hamiltonian $\hat{H}$

$$\hat{H}\underline{b}^{(k)}=w_k\underline{b}^{(k)}. \quad (3)$$

The quantity $r_k(t)$ is the statistical weight of the adiabatic dressed-state vector $\underline{b}^{(k)}(t)$ in the (bare) state vector $\underline{a}(t)$. According to the adiabatic theorem [18], $r_k(t)\equiv r_k(-\infty)$: If the system is in a (bare) eigenstate of $\hat{H}$ before the action of the laser pulse it will be passed into the (dressed) eigenstate of $\hat{H}$ that derives from the bare eigenstate by continuity when the interaction with the laser pulse is being switched on. When the evolution is adiabatic no transitions between the adiabatic states take place and their populations are preserved.

We obtain the following equation for the eigenvalues (quasienergies) w_k, $k=1,2,\ldots,6$ from Eq. (3):

$$g_{11}g_{22}-g_{12}g_{21}=0, \quad (4)$$

where

$$g_{11}=\frac{w\Omega_{14}}{\Omega_{13}}+\frac{2(w+\varepsilon_4)\Omega_{25}}{\Omega_{25}\Omega_{14}-\Omega_{15}\Omega_{24}}\left[\frac{\Omega_{13}}{2}-\frac{2w(w+\varepsilon_1)}{\Omega_{13}}\right],$$

$$g_{12}=\frac{(w+\varepsilon_6)\Omega_{24}}{\Omega_{26}}-\frac{2(w+\varepsilon_4)\Omega_{15}}{\Omega_{25}\Omega_{14}-\Omega_{15}\Omega_{24}}\times\left[\frac{\Omega_{26}}{2}-\frac{2(w+\varepsilon_2)(w+\varepsilon_6)}{\Omega_{26}}\right],$$

$$g_{21}=\frac{w\Omega_{15}}{\Omega_{13}}-\frac{2(w+\varepsilon_5)\Omega_{24}}{\Omega_{25}\Omega_{14}-\Omega_{15}\Omega_{24}}\left[\frac{\Omega_{13}}{2}-\frac{2w(w+\varepsilon_1)}{\Omega_{13}}\right],$$

$$g_{22}=\frac{(w+\varepsilon_6)\Omega_{25}}{\Omega_{26}}+\frac{2(w+\varepsilon_5)\Omega_{14}}{\Omega_{25}\Omega_{14}-\Omega_{15}\Omega_{24}}\times\left[\frac{\Omega_{26}}{2}-\frac{2(w+\varepsilon_2)(w+\varepsilon_6)}{\Omega_{26}}\right].$$

Ω_{ij} $(i,j=1,\ldots,6)$ are Rabi frequencies for corresponding transitions and detunings ε_i $(i=1,\ldots,6)$ are defined in Eq. (2).

The time dependencies of solutions of Eq. (4), the eigenvalues (quasienergies) w_k, $k=1,2,\ldots,6$, are shown in Fig. 5 for the linearly chirped laser pulse with a Gaussian envelope. The diabatic lines (the eigenvalues w_k, $k=1,2,\ldots,6$, versus time in the absence of the laser field when $\Omega_{ij}\rightarrow 0$) are depicted in Fig. 5 as dashed lines.

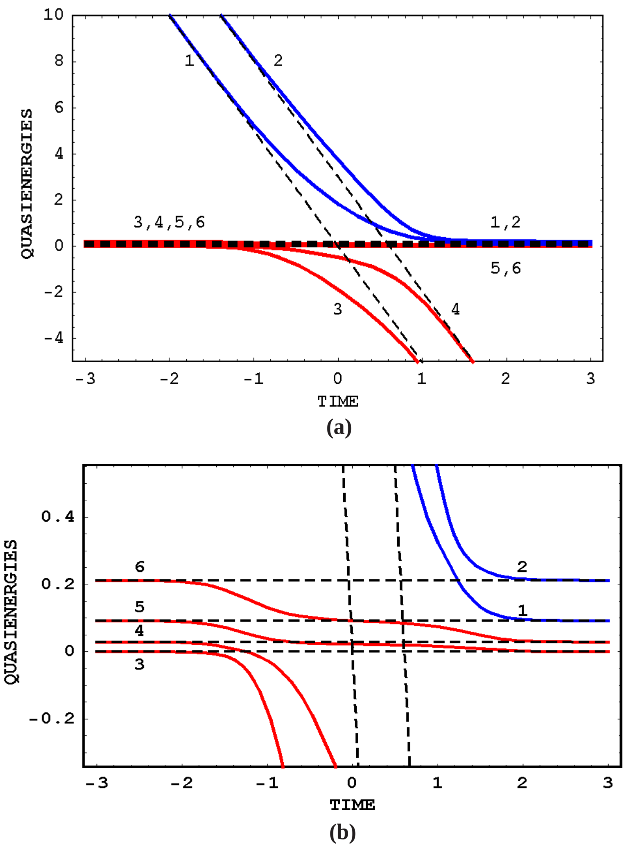


FIG. 5. (Color online) (a) Time dependence of the quasienergies of the hyperfine levels (solid lines). Diabatic lines are shown by dashed lines. Time is normalized by the duration of the laser pulse. The parameters applied are the duration of the pulse (intensity) $2\tau_p = 2$ ns; the peak Rabi frequency is between $\Omega_R = 1.73$ GHz and $\Omega_R = 3$ GHz for the allowed transitions between the two ground and four excited states, the negative (from blue to red) chirp speed is $b = -3$ GHz/ns. (b) A magnified picture for small values of the quasienergies.

We solve Eq. (3) to find the components $b_i^{(k)}$ of the dressed state $\underline{b}^{(k)}(t)$ for corresponding quasienergy w_k. The solutions for the corresponding probability amplitudes are

$$b_i^{(k)} = C_i(w_k)/\sqrt{N(w_k)} \quad (i = 1, \dots, 6), \tag{5}$$

where

$$C_1 = \frac{2w_k}{\Omega_1} C_3,$$

$$C_2 = -\frac{2(\varepsilon_6 + w_k)}{\Omega_6^*} C_6,$$

$$C_3 = \frac{(\varepsilon_6 + w_k)\Omega_4^*}{\Omega_6^*} - \frac{2(\varepsilon_4 + w_k)}{\Omega_5\Omega_2 - \Omega_3\Omega_4}\Omega_3 \times \left[\frac{\Omega_6}{2} - \frac{2(\varepsilon_2 + w_k)(\varepsilon_6 + w_k)}{\Omega_6^*}\right],$$

$$C_4 = -C_3 \frac{2\Omega_5}{\Omega_5\Omega_2 - \Omega_3\Omega_4}\left[\frac{\Omega_1}{2} - \frac{2w_k(\varepsilon_1 + w_k)}{\Omega_1^*}\right] - C_6 \frac{2\Omega_3}{\Omega_5\Omega_2 - \Omega_3\Omega_4}\left[\frac{\Omega_6}{2} - \frac{2(\varepsilon_2 + w_k)(\varepsilon_6 + w_k)}{\Omega_6^*}\right],$$

$$C_5 = C_3 \frac{2\Omega_4}{\Omega_5\Omega_2 - \Omega_3\Omega_4}\left[\frac{\Omega_1}{2} - \frac{2w_k(\varepsilon_1 + w_k)}{\Omega_1^*}\right] + C_6 \frac{2\Omega_2}{\Omega_5\Omega_2 - \Omega_3\Omega_4}\left[\frac{\Omega_6}{2} - \frac{2(\varepsilon_2 + w_k)(\varepsilon_6 + w_k)}{\Omega_6^*}\right],$$

$$C_6 = \frac{w_k\Omega_2^*}{\Omega_1^*} + \frac{2(\varepsilon_4 + w_k)}{\Omega_5\Omega_2 - \Omega_3\Omega_4}\Omega_5\left[\frac{\Omega_1}{2} - \frac{2w_k(\varepsilon_1 + w_k)}{\Omega_1^*}\right].$$

The normalization factor is $N(w_k) = \Sigma |C_i|^2$.

We construct the dressed state $\underline{b}^{(k)}(t)$ by selecting an eigenvalue w_k from the solution of Eq. (4) which tends to the diabatic line that corresponds to a given initial conditions for the atom before interaction with the laser pulse at $t \to -\infty$ $(\Omega_{ij} \to 0)$. For example, if the atom was in the state $|1\rangle$ initially, we have to choose a solution w_1, which tends to the diabatic line 1 in Fig. 5 at $t \to -\infty$, corresponding to the quasienergy of the bare state $|1\rangle$. Using w_1, we obtain the dressed-state vector $\underline{b}^{(1)}(t)$ from Eq. (5). The solution of Eq. (4), w_2, that tends to the diabatic line 2 in Fig. 5, is the quasienergy of the atom that initially was in the ground state $|2\rangle$. In the case when the population of the atom is distributed with some statistical weights between the two (bare) ground states initially, we have to construct combination of the dressed states $\underline{b}^{(1)}(t)$ and $\underline{b}^{(2)}(t)$ with the same statistical weights as for the bare states. As it follows from Fig. 5, as a result of the interaction with a laser pulse having a negative frequency chirp (frequency is diminishing during the pulse, blue to red chirp), the quasienergies w_1 and w_2 evolve into the diabatic lines corresponding to the excited states $|5\rangle$ and $|6\rangle$, respectively, which are the uppermost levels for transitions from the ground states $|1\rangle$ and $|2\rangle$, respectively (compare with the results of Ref. [16]).

Mathematically, adiabatic evolution requires that coupling between the adiabatic states is negligibly small compared with the difference between their quasienergies. In the simplest case of a two-level system, the adiabatic condition reads [13]

$$\left|\varepsilon \frac{d}{dt}\Omega - \Omega \frac{d}{dt}\varepsilon\right| \ll 2(\Omega^2 + \varepsilon^2)^{3/2},$$

where Ω and ε are the Rabi frequency and detuning from the resonance for transitions in the two-level system. As it follows from this condition, the adiabatic evolution requires a smooth laser pulse, relatively long interaction times, and large Rabi frequencies. While the explicit form of the adiabatic condition is relatively simple in the case of a two-level system it is much more complicated in the case of a multilevel system under consideration in this paper.

The population dynamics of the hyperfine levels of the ^{85}Rb atom model calculated using the numerical simulation

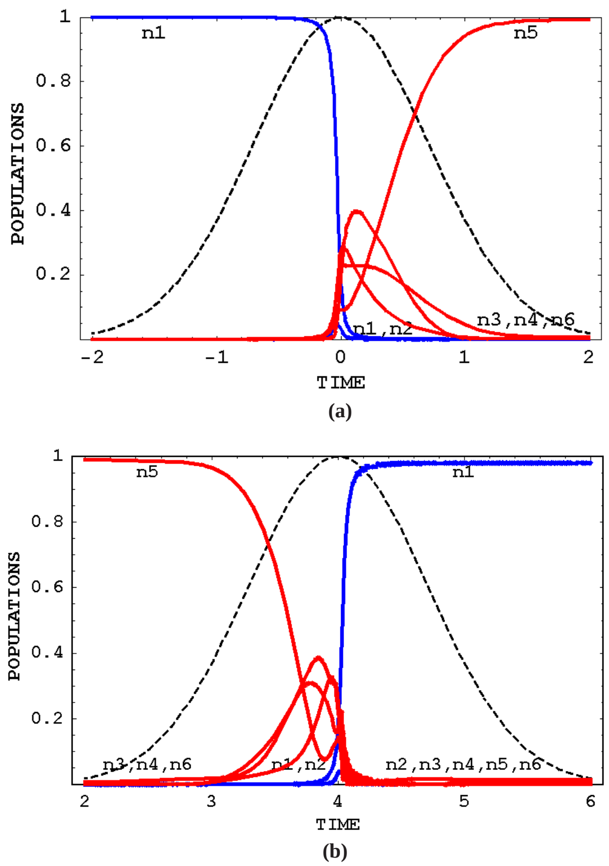


FIG. 6. (Color online) Numerical simulation: The population dynamics in the case of long pulses when all the hyperfine levels are resolved. Time is normalized by the duration of the laser pulse. The parameters applied are the duration of the pulse (intensity) $2\tau_p = 50$ ns; the peak Rabi frequency is between $\Omega_R = 2.61$ GHz and $\Omega_R = 4.5$ GHz for the allowed transitions between the two ground and four excited states. The dashed line is the pump pulse intensity profile. (a) During the first pulse where the negative (from blue to red) chirp speed is $b = -1/2$ GHz/ns. (b) During the second pulse where the positive (from red to blue) chirp speed is $b = 1/2$ GHz/ns.

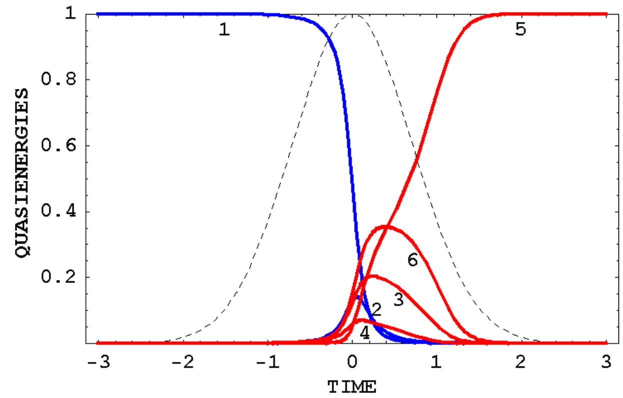


FIG. 7. (Color online) The population dynamics during the first pulse calculated in the dressed-state picture. The parameters are the same as in Fig. 6(a) except for the Rabi frequency, which is 15 times smaller.

of the Schrödinger equation is shown in Fig. 6 for the case of the atom in the ground state $|1\rangle$ initially. As it follows from this figure, all the atomic population is transferred into the excited state $|5\rangle$ as a result of action of the frequency-chirped laser pulse with the blue to red direction of the frequency chirp [see Fig. 6(a)]. One has to apply a laser pulse with the opposite (red to blue) direction of the frequency chirp to move the atomic population back to the ground state $|1\rangle$ [see Fig. 6(b)]. The results of the dressed-states analysis [see Eqs. (4) and (5)] are depicted in Fig. 7 and show good agreement with the results based on the numerical simulation (see Fig. 6). This agreement shows that the parameters applied during the numerical simulation are in the range of fulfillment of the adiabatic conditions for interaction of the frequency-chirped laser pulse with the Rb model system.

V. CONCLUSIONS

In conclusion, the results of the analysis of population transfer in the ^{85}Rb atom in the field of a sequence of frequency-chirped laser pulses are presented in this paper. We have shown that effective population transfer between two ground hyperfine levels of the ^{85}Rb atom and the manifold of excited hyperfine levels may be achieved using frequency-chirped laser pulses of different durations: relatively short (broadband pumping), relatively long (narrow band pumping), and also for laser pulses long enough for all working hyperfine levels in both ground and excited levels manifolds to be resolved. In all regimes mentioned, the frequency of the laser pulse must be chirped through the resonance with all allowed transitions between the two ground and four excited hyperfine working levels of the ^{85}Rb atom to obtain the effective population transfer. Dependence of the population transfer to individual excited states on the initial coherence between the two ground states is demonstrated in the regime of the broadband pumping when the frequency spectrum of the laser pulse envelope exceeds the frequency interval between the ground states. For example, destructive quantum interference may suppress the transfer of the atomic population to excited hyperfine levels (levels $|4\rangle$ and $|5\rangle$ in the ^{85}Rb atom, see Fig. 1), which are common for transitions from both ground states. On the contrary, there is enhancement of the population transfer to these excited states under conditions of constructive interference. It is important to note, however, that independently on the initial coherence between the ground states, population transfer takes place from the two ground states to the manifold of excited states and back to the ground states as a result of interaction with the frequency-chirped laser pulses in the broadband pumping regime.

We have shown that effective population transfer takes place for a narrow band pumping, too, when the ground hyperfine levels are resolved but the excited ones are not. Analysis shows that the effective population transfer takes place in both cases when the atomic population is optically pumped into one of the ground states or it is located in both ground states initially. In the latter case, the efficiency of the population transfer does not depend on the initial coherence between the ground states.

For the case of frequency-chirped laser pulses long enough for all hyperfine levels to be resolved, we have performed dressed-states analysis and a numerical simulation of the Schrödinger equation to show that in a Rb atom model, population transfer from the two ground states to excited ones and back to the ground states takes place when conditions for the AP regime of interaction are fulfilled.

As it follows from the above consideration, the interaction of a frequency-chirped laser pulse with a ^{85}Rb atomic system in the ground states initially results in excitation of the atom in the case of a sufficiently high Rabi frequency of the pulse: the all atomic population becomes distributed among the excited states with no population in the ground states just similar to the interaction of a frequency-chirped laser pulse with an effective two-level system in the AP regime of interaction. Second laser pulse(s) applied to the atom, with a time delay much shorter than the relaxation times of the atom, will transfer all atomic population into the ground states. This may be used for the efficient coherent manipulation of beams of multilevel atoms, such as Rb atoms: the atoms will receive mechanical momentum equal to $2\hbar k_p$ (with k_p being the wave number of the laser pulse) after interaction with a pair of frequency-chirped laser pulses, which propagate in opposite directions. Repetition of this cycle of excitation and deexcitation of the atoms by a sequence of counterpropagating laser pulses will result in the effective coherent transfer of the mechanical momentum from the laser field to Rb atoms and the effective manipulation of the atomic beam. An important advantage of this system is that relatively large mechanical momentum may be transferred to the atomic beam from the laser field due to a large number of the laser pulses interacting with the atoms. This number is defined by the repetition rate of the laser pulse source and the time of flight of the atoms through the interaction region.

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