Comment on "Realization of a bipolar atomic Šolc filter in the cavity-QED microlaser"

M. A. Bouchene
Laboratoire de Collisions Agrégats Réactivité, CNRS UMR 5589, IRSAMC, Université Paul Sabatier, 118 Route de Narbonne, F-31062 Toulouse Cedex 4, France

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Seo *et al*. [Phys. Rev. A **81**, 053824 (2010)] have studied the interaction of a TEM_{10} mode in a cavity-QED microlaser. They claim that this scheme represents an experimental realization of an atomic Šolc filter proposed by Hong *et al*. [Opt. Express **17**, 15455 (2009)]. Even if the regime of interaction considered is nonperturbative, we point out that this statement is valid only for a small range of intensity and is misleading in the general case of an arbitrarily strong field where the effect is related to the nonadiabatic jump phenomenon [G. S. Vasilev and N. V. Vitanov, Phys. Rev. A **73**, 023416 (2006)].

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Seo *et al*. [1] report the results of an experimental study on the interaction of an initially inverted two-level atom that crosses a TEM_{10} mode interacting with the cavity field in a strong-coupling regime. The authors interpret their results by invoking a flat bipolar coupling model connecting this experiment to the atomic Šolc filter described in Ref. [2]. Next, to develop our arguments, we point out that the interaction of a TEM_{10} mode with a two-level system with arbitrary detuning and intensity has already been studied intensely in Ref. [3] [shape envelope equation (8c)]. It is then natural to explain the underlying physics at the highlight of this Reference. Other envelope shapes have been studied in Refs. [4,5] for coherent control purposes. Also, we want to emphasize that, in this Comment, our criticism neither concerns the theoretical model developed for the atomic Šolc filter in Ref. [2] nor the numerical simulation in Ref. [1] [Figs. 3(d) and 4 account for the exact shape of the pulses and, thus, are not questionable]. Our criticism concerns the interpretation of the present experimental scheme (with a TEM_{10} pulse) as an experimental realization of the atomic Šolc filter (as treated in Ref. [2]) in the general case of an arbitrary intense field. This point was missed by the authors in their discussion. Indeed, the authors have experimentally limited their investigation to relatively moderate power. The regime is nonperturbative, and the analogy with the atomic Šolc filter is justified ($m \simeq 1$), but the exact picture of the dynamics of the phenomenon for larger field intensities goes far beyond the crude approximation of a flat coupling model, excluding the extent of the results of Ref. [2] to this regime. What Ref. [3] shows (through Fig. 6, for instance), is that no matching conditions are generally required in the case of a TEM_{10} mode: In the presence of detuning, an almost complete inversion population is obtained if the Rabi frequency is high enough. Since no matching conditions are required and inversion can be obtained for (arbitrary) strong fields with detuning, the analogy with the Šolc filter is not pertinent in the general case. Except for the small experimental zone explored, the atomic Šolc filter cannot be approached by smooth varying pulses with adapted detuning and Rabi frequencies.

The realistic case with a smooth varying envelope (Gaussian, sech, etc.) differs qualitatively and quantitatively with the particular case where the electric field switches suddenly and not smoothly (at the beginning and the end of the pulse and not only at the period of sign change). In the latter, the dynamics of the systems is dominated by Rabi oscillations, whereas, in the former, a subtle interplay between adiabatic following and Rabi oscillations occurs.

Next, we briefly explain the concrete case described in detail in Ref. [3] and then consider some other fundamental differences with a flat bipolar coupling model. When taking the analytical form of the pulse into account, the dynamics is obtained in the adiabatic basis, which is defined by

$$\begin{pmatrix} |+,n\rangle \\ |-,n\rangle \end{pmatrix} = \begin{bmatrix} \cos[\theta_n(t)/2] & \sin[\theta_n(t)/2] \\ -\sin[\theta_n(t)/2] & \cos[\theta_n(t)/2] \end{bmatrix} \begin{pmatrix} |g,n\rangle \\ |e,n-1\rangle \end{pmatrix}. \quad (1)$$

Here, $|\pm,n\rangle$ are the dressed states and $|g,n\rangle$, $|e,n-1\rangle$ are the bare states. The mixing angle θ_n is defined by $\tan[\theta_n(t)] = 2(\eta/\delta)g(t)/g_0$ with η, δ, $g(t)$ and g_0 as the normalized Rabi frequency, normalized detuning, TEM_{10} atom-cavity coupling, and flat bipolar equivalent model atom-cavity coupling, respectively. These quantities are defined, such as in Ref. [1]. Writing the Schrödinger equation for the system, the amplitudes $\alpha_{\pm,n}$ of states $|\pm,n\rangle$ obey the following equation:

$$\partial_T \begin{pmatrix} \alpha_{-,n} \\ \alpha_{+,n} \end{pmatrix} = \begin{pmatrix} i\Omega_n(T)/2 & -\partial_T\theta_n \\ \partial_T\theta_n & -i\Omega_n(T)2 \end{pmatrix} \begin{pmatrix} \alpha_{-n} \\ \alpha_{+n} \end{pmatrix}(T), \quad (2)$$

where $\partial_T = \partial/\partial T$, $T = t/\tau$, and $\Omega_n = [\delta^2 + 4ng^2(T)]^{1/2}$ is the dimensionless generalized Rabi frequency. The diagonal terms $\pm\Omega_n/2$ represent the energy levels of the dressed states, and $\partial_T\theta_n$ is the nonadiabatic coupling. Ω_n gives the instantaneous separation of the dressed levels becoming δ as T approaches $\pm\infty$ [with $|-,n\rangle(\pm\infty) = |g,n\rangle$ and $|+,n\rangle(\pm\infty) = |e,n-1\rangle$]. For pulses such that $|\partial_T\theta_n/\Omega_n| \ll 1$, the evolution can be considered to be adiabatic, the dressed levels experience a transient light shift during the action of the driving pulse, and no population is transferred asymptotically. In the opposite case, there can be nonadiabatic transitions to the level $|-,n\rangle$ that result in an asymptotic population in atomic state $|g\rangle$. The key point here is that the transfer strongly depends on the shape of the pulse through $\partial_T\theta_n$, and this dependence can lead asymptotically to a complete population inversion. This is the case for asymmetric pulses (such as a TEM_{10} pulse or other forms studied in Refs. [3,4]). Indeed, it has been shown that the coupling approximates to a π-area δ-like function in the limit of strong driving ($\eta/\delta \gg 1$) explaining the surprising efficiency of such a nonresonant process (see Fig. 3 in Ref. [3] or Fig. 1 in Ref. [4]). The very abrupt transfer obtained in this case justifies the nonadiabatic jump (NAJ) denomination for this phenomenon.

So, the NAJ phenomenon leads to population inversion, such as the atomic Šolc filter, but the phenomena underlying the inversion process are different.

The dynamics can also be explained in a Bloch picture. For $\eta/\delta \gg 1$, the Bloch vector evolves *adiabatically* from the north pole ($t \to -\infty$) to the equator ($t = 0$) and from the equator to the south pole ($t \to +\infty$). Around $t = 0$, NAJ takes place ensuring that the Bloch vector does not evolve, whereas, the rotating vector changes to the opposite position. Paradoxically, the role of NAJ is to invert the population in adiabatic states, whereas, in bare states it is to maintain the Bloch vector in its position during the electric sign change step. Thus, the population inversion is the consequence of the combined role of adiabatic switching and NAJ. However, because of NAJ, the sign change can be realized with smooth pulses and not abrupt ones. So, the TEM_{10} case is not just a form close to the ideal situation of an atomic Šolc filter but represents both the ideal and the realistic cases.

In the experiment performed by the authors, the field parameters were such that adiabatic following could not take place and corresponded to the zone around the small lobe ($\Delta \simeq \sqrt{2}$ and $\sqrt{\Omega_0} \simeq 2$) of Fig. 6 in Ref. [3]. In this situation, the Bloch vector evolves from the north pole to the equator and from the equator to the south pole under Rabi oscillations as the authors state. However, and as mentioned above, this interpretation cannot be extrapolated for the case of stronger pulses. Moreover, by comparing Fig. 6 to Figs. 4 and 5 of Ref. [3], an important effect appears: The change in the shape of the envelope also changes the domain of the parameters for which the process is efficient. Thus, the interpretation of concrete cases cannot be reduced to a bipolar flat model without missing important effects and has to be performed within the framework of Ref. [3].

Invoking a flat bipolar model also leads to many misleading consequences:

(1) The inversion population in the NAJ model (and, thus, for the TEM_{10} mode), confirmed by numerical simulations (Fig. 6 of Ref. [3]) is robust with excitation parameters (Rabi frequency, detuning, and pulse duration). This is in contrast with Eq. (1) in Ref. [1] that leads to the incorrect statement that the population in the excited state oscillates between 0 and 1 when the interaction time is tuned, for instance. Compare, for instance, Fig. 6 of Ref. [3] and the insets in Fig. 3 of Ref. [1]. This strengthens the importance of the fact that no matching conditions are required here to obtain the population inversion, whereas, in the atomic Šolc filter, adapted Rabi oscillations are required to achieve this purpose.

(2) The situation described in the papers [3,4] is considered for pulses with duration and detuning such that $\delta \gg 1$. It ensures that the frequency components at resonance are negligible. For instance, for a Gaussian pulse and $\delta = 8$, we have $S \propto e^{-\delta^2} \simeq e^{-64} \simeq 10^{-28}$ for $\delta \gg 1$ (S is the intensity spectrum at atomic resonance). In this case, the presence of residual resonant frequencies cannot be invoked to explain the population inversion. Even if the coupling $\dot{\theta}_n$ reveals a singular behavior (δ-like function with π area), the electric field is smooth. This is in contrast with the atomic Šolc filter case for which the spectrum exhibits some long-range tails whose influence cannot be neglected in the strong-field regime.

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