Optically Controllable Photonic Structures with Zero Absorption

Chris O'Brien* and Olga Kocharovskaya
Department of Physics and Astronomy and Institute for Quantum Studies, Texas A&M University, College Station, Texas 77843-4242, USA

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We show the possibility to periodically modulate the refractive index in a homogeneous resonant atomic medium in space or/and time while simultaneously keeping vanishing absorption or gain. Such modulation is based on periodic resonant enhancement of the refractive index, controlled by an external optical field, and opens the way to produce coherently controllable photonic structures. We suggest the possible implementation of the proposed scheme in rare-earth doped crystals with excited state absorption.

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One, two, or three-dimensional periodic heterostructures made of two dielectric materials with different refractive indices, such as distributed Bragg reflectors (DBRs), holey fibers, or photonic crystals find many applications, including reflective coatings, distributed feedback lasers, and optical cavities. Different technologies such as photolithography, etching, drilling, and self-assembling are used for construction of such structures.

We suggest a method to produce transparent photonic structures in a homogeneous resonant atomic media, such as dielectrics with homogeneously distributed impurities, atomic, or molecular gases, simply by illuminating these materials with standing waves of a laser field. Such optically produced photonic structures could easily be controlled (including switching on or off, changing amplitude and period of modulation) and would be highly selective in frequency, naturally limited by the width of the optical resonance.

Refractive index (RI) is strongly enhanced near atomic resonances. However, that enhancement is accompanied by enhancement of absorption. Namely, when the maximal contribution from the atomic resonance to the RI is reached, the contribution to the absorption is on the same order which prevents the usage of obtained RI. There have been several proposals on how to resonantly enhance the refractive index while at the same time eliminating resonant absorption. One approach is based on interference effects in multilevel atomic systems driven by coherent resonant fields [1–5]. Another suggestion is to compensate absorption with resonant gain from an inverted transition [6]. Such a situation could be realized either in a mixture of two two-level atomic species, or in a single atomic species possessing simultaneously both noninverted and inverted transitions with slightly shifted frequencies [7]. Proof of principle experiments were done in hot Rb vapors in which enhancement of the refractive index $\Delta n \sim 10^{-4}$ was achieved under negligible absorption [8,9]. An enhancement up to the value $\Delta n \sim 10^{-2}$ is expected with an increase of density to $N = 6 \times 10^{16}$ cm^{-3}. The further increase of the refractive index in room-temperature gases is not feasible due collisional broadening becoming the dominant contribution to the linewidth. Much higher resonant additions to the background index are anticipated in transition element doped crystals due to the essentially higher density of the ions which does not in general result in proportional line broadening [7,10,11].

In all of these proposals the RI was uniform in space. Moreover, an enhancement of the RI with vanishing resonant absorption was achieved only at a particular detuning of the probe field from atomic resonance and was accompanied by either absorption or gain at the neighboring detunings. Thus, none of those proposals were suitable for achieving spatial modulation of the refractive index with zero absorption. Our proposal is based on spatial modulation of the energy of a populated intermediate state in a nearly degenerate ladder configuration via the ac-Stark effect in a standing wave field which results in a spatially dependent detuning leading to a periodic resonant increase and decrease of the refractive index in space while simultaneously keeping transparency of the medium.

Consider the interaction of a probe field with a medium of three level atoms in a ladder configuration such that the probe field interacts with both transitions as illustrated in the inset of Fig. 1. The transition frequencies ω_{21} and ω_{32} are close to each other so that the probe field with frequency ω_p interacts simultaneously with both transitions and for a weak probe Rabi frequency $\Omega_p \ll \gamma_{21}, \gamma_{32}$ the susceptibility is defined as the sum of the susceptibilities of two two-level transitions:

$$\chi = \frac{3N\lambda^3}{8\pi^2}\left[\frac{\gamma_{21}^{\mathrm{rad}}(\rho_1 - \rho_2)}{\delta_{21} - i\gamma_{21}} + \frac{\gamma_{32}^{\mathrm{rad}}(\rho_2 - \rho_3)}{\delta_{32} - i\gamma_{32}}\right]. \quad (1)$$

Here N is the atomic density, the detunings are defined as $\delta_{21} = \omega_{21} - \omega_p$ and $\delta_{32} = \omega_{32} - \omega_p$, λ is the probe field wavelength in the medium, $\gamma_{ij}^{\mathrm{rad}}$ is the radiative decay rate for the i to j transition, γ_{ij} is the total decoherence rate, and ρ_i is the population in the ith energy level. We assume that

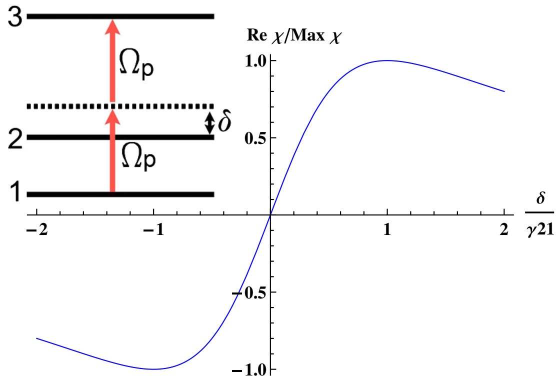

FIG. 1 (color online). Real part of the susceptibility as a function of the level shift δ. Note that the imaginary part is identically zero. Inset: the energy level diagram for the corresponding three-level scheme.

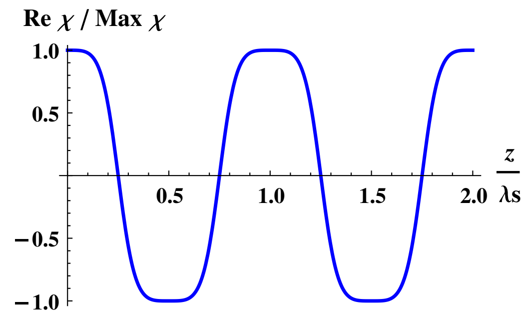

FIG. 2 (color online). Real part of the susceptibility plotted as a function of position along the optical axis.

the amplitudes of both transitions are matched but of opposite sign:

$$\gamma_{21}^{\rm rad}(\rho_1 - \rho_2) = -\gamma_{32}^{\rm rad}(\rho_2 - \rho_3), \qquad (2)$$

which means that one of the two transitions is inverted. Let it be transition 2-1, i.e. $\rho_2 - \rho_1 > 0$. We also assume the widths of the transitions are equal $\gamma_{21} = \gamma_{32}$ and the probe field is tuned to two-photon resonance, i.e. $\omega_p = \omega_{31}/2$. Thus for arbitrary position of level 2 the blue detuning of the probe field from one of two two-level transitions is equal to the red detuning from the other, i.e., $\delta_{32} = -\delta_{21} = \delta$, leading to the remarkable property that gain at one transition and absorption at another one cancel each other while the real part of susceptibility is doubled. So, the susceptibility is purely real:

$$\chi = \frac{3N\lambda^3\gamma_{21}^{\rm rad}(\rho_1 - \rho_2)}{8\pi^2}\,\frac{2\delta}{\delta^2 + \gamma_{21}^2}. \qquad (3)$$

It means that the probe field neither experiences absorption nor gain independently of level 2's energy, i.e., for arbitrary values of δ. At the same time the resonant susceptibility varies from the minimum to the maximum value as δ is shifted from $-\gamma$ to γ as shown in Fig. 1. If the energy of the intermediate level is modulated in space along the direction of propagation of the probe field, the refractive index is also modulated. Such spatial modulation can be produced along the optical axis via the ac-Stark shift. A control laser field $E_s\cos(\omega_s t)$ applied at the 0-2 transition adjacent to the 1-2 transition and far detuned from this transition $\Delta_s = \omega_s - \omega_{20} \gg \gamma_{20}$ would result in a splitting of the intermediate state 2 into two ac-Stark sublevels shifted in frequency by $-|\Omega_s|^2/\Delta_s$ and $\Delta_s + |\Omega_s|^2/\Delta_s$, respectively, where Ω_s is the associated Rabi frequency. The probe field is far out of resonance with the transitions from the second Stark sublevel from both level 1 and level 3 and, therefore its interaction with these transitions is negligible while the first Stark sublevel is slightly shifted from the original level 2 and strongly interacts with the probe field. In other words, the susceptibility at each transition (2-1 or 2-3), which in general consists of two terms associated with the one-photon and two-photon resonances is reduced to the one-photon contribution and has the same form as given by Eq. (3), just with shifted transition frequencies.

If the control field represents itself as a standing wave such that the Rabi frequency is a function of position inside the medium, $\Omega_s(z) = \Omega_s\cos(k_s z)$, then the ac-Stark shift of level 2 is given by:

$$\Delta E = -\frac{\hbar|\Omega_s|^2}{2\Delta_s} - \frac{\hbar|\Omega_s|^2}{2\Delta_s}\cos(2k_s z). \qquad (4)$$

Thus it consists of a constant shift, $|\Omega_s|^2/2\Delta_s$, and a sinusoidal modulation, $(|\Omega_s|^2/2\Delta_s)\cos(2k_s z)$. If the difference between the atomic transition frequencies $\omega_{32} - \omega_{21}$ is chosen to be equal to $-|\Omega_s|^2/\Delta_s$ then the susceptibility is described by Eq. (3) with $\delta = (|\Omega_s|^2/2\Delta_s)\cos(4\pi z/\lambda_s)$ (where λ_s is the wavelength of the control field in the medium). Hence the refractive index will be modulated symmetrically with respect to its background value as shown in Fig. 2. The spacial period $\lambda_s/2$ is defined by the wavelength, while the modulation depth $-|\Omega_s|^2/\Delta_s$ is defined by the Rabi frequency of the modulating field Ω_s. To provide the maximum amplitude of refractive index modulation the Rabi frequency of the control field should meet the condition $\Omega_s^2 = 2\gamma\Delta_s$.

With a strong enough index variation a transparent for a particular frequency 1-D photonic crystal can be created with properties that are optically controlled. Similarly a 2D or 3D photonic structure can be produced by application of 2 or 3 orthogonal modulating control fields. Even for index variations much smaller than the background RI the medium will behave as a distributed Bragg reflector if $\lambda_s \simeq \lambda_p$ specifically, when the wavelength mismatch is within the width of the Bragg band gap, $\lambda_{\rm s} - \lambda_{\rm p} < \lambda_s\Delta n/(\pi n_{\rm bg})$. Since the medium remains transparent, many periods of spatial RI structures can be used as needed to achieve the required reflection coefficient.

When the probe field is detuned from two-photon resonance with 1-3 transition it will experience either gain or absorption. The question arises if such gain may result in the building up of a spontaneously amplified field emptying the inverted transition and limiting the propagation length of the probe field in the medium with periodic refractive index. Fortunately, this is not the case. Indeed, since the position of the intermediate level is periodically modulated in space, then a detuned probe field experiences periodically interchanging regions of gain and absorption suppressing the development of such an instability as can be seen in Fig. 3. In fact averaging the absorption over a wavelength λ_s shows that the medium is effectively transparent even when the probe field is detuned from resonance.

The simple model of a ladder system previously discussed assumed the existence of two transitions possessing equal linewidths, equal products of transition strength and population difference, and nearly degenerate (on the scale of the linewidth) frequencies. It is difficult if not impossible to meet these conditions in a real atomic system. However, it is possible to construct an effective ladder system whose upper transition has controllable parameters which could be optically tuned to satisfy these conditions.

It can be accomplished by adding to the original simple ladder system along with the modulating control field E_s coupled to an adjacent transition 0-2 (as discussed above) a second control field E_c coupling the excited state 3 to an additional unpopulated level 4 as shown in Fig. 4. This second far-detuned control field ($\Delta_c \gg \gamma_{32}^{\rm rad}$, Ω_c where $\Delta_c = \omega_{43} - \omega_c$ and Ω_c is the control field Rabi frequency) is chosen to satisfy approximately the two-photon resonance condition: $\omega_c - \omega_p = \omega_{42}$, forming together with the probe field a far-detuned lambda scheme. A strong far-detuned field results in an ac-Stark splitting of level 3 and the response to the probe field consists of two terms representing one-photon (upper Stark sublevel) and two-photon (lower Stark sublevel) contributions in the same way as previously discussed. But now it is the two-photon contribution which plays a dominant role due to the two-photon resonance condition [7,12].

As a result, the total five level system under the formulated above conditions is reduced to an effective three-level ladder system with the lower transition 1-2'' and the upper transition 2''-3'. Its susceptibility takes the form:

$$\chi_{\rm res} = \frac{3N\lambda^3}{8\pi^2}\left\{\frac{\gamma_{21}^{\rm rad}p/(2-p)}{\delta_p + \frac{\Omega_s^2}{2\Delta_s} + \delta - i\gamma_{21}} + \frac{\xi\gamma_{32}^{\rm rad}/[(2-p)(1+2\xi)]}{\delta_p - \omega_{32} + \omega_{21} - \frac{\Omega_s^2}{2\Delta_s} - \delta + \Delta_c(1+\xi-\xi^2) - i[\gamma_{42}(1-\xi)+\gamma_{32}\xi]}\right\}. \quad (5)$$

Where we assume incoherent pumping (not shown in Fig. 4) which provides the necessary population inversion, represented by the pumping factor $p = (\rho_2 - \rho_1)/\rho_2$. We also assume level 3 is empty and introduce a control field parameter $\xi = |\Omega_c|^2/\Delta_c^2$, as well as the one-photon detuning $\delta_p = \omega_{21} - \omega_p$. Now the parameters of the effective upper 2''-3' and lower 1-2'' transitions defined by the control fields can easily be matched.

We choose $\Omega_s = \sqrt{2\gamma_{21}\Delta_s}$ to provide the maximum range of refractive index modulation. Matching the linewidth of 3'-2'' transition to that of 2''-1 defines the control field parameter ξ:

$$\xi = \frac{\gamma_{21} - \gamma_{42}}{\gamma_{32} - \gamma_{42}}. \quad (6)$$

It implies a larger linewidth of the upper 2''-3' transition as compared to the lower transition 1-2'', $\gamma_{32} > \gamma_{21}$, and relatively slow decay of the coherence at the 4-2''

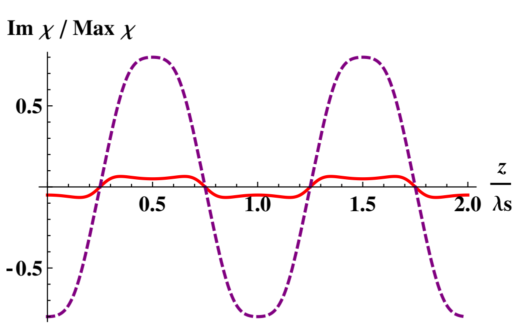

FIG. 3 (color online). Imaginary part of the susceptibility for a probe field detuned from resonance by $\gamma_{21}/20$ (solid line) and γ_{21} (dashed line) plotted as a function of position along the optical axis.

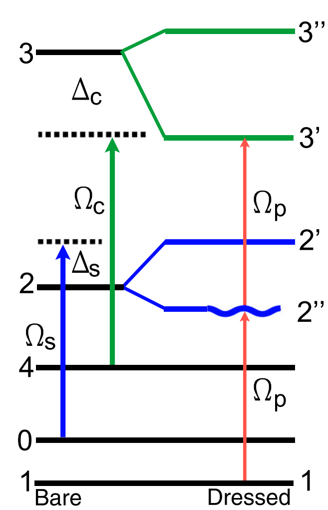

FIG. 4 (color online). Energy level diagram for the 5-level system coupled with two control fields Ω_s and Ω_c leading in ac-Stark splitting of levels 2 and 3 and resulting in an effective ladder system 1-2''-3' in the dressed state basis.

transition: $\gamma_{42} < \gamma_{32}$, γ_{21}. Matching the amplitudes defines the pump parameter as:

$$p = \frac{\gamma_{32}^{\rm rad}}{\gamma_{21}^{\rm rad}} \frac{\xi}{1 + 2\xi}. \quad (7)$$

We take the probe field to be resonant with the dressed transition 2"-1 such that $\delta_p = -\Omega_s^2/2\Delta_s$. Then matching the frequencies of the transitions defines the required detuning of the control field Δ_c:

$$\Delta_c = \frac{\omega_{32} - \omega_{21} + 2\gamma_{21}}{1 + \xi - \xi^2}. \quad (8)$$

This implies that Δ_c will be on the same order as $\omega_{32} - \omega_{21}$. Since $|\Omega_c| = \sqrt{\xi}\Delta_c$ and $\Delta_c \approx \omega_{32} - \omega_{21}$, it is important to have 1-2 and 2-3 transitions with close frequencies in order to reduce the required control field intensity. Under the above conditions the susceptibility given by Eq. (5) takes the same form as in Eq. (3). Thus, it becomes possible to realize resonant modulation of the refractive index with zero absorption or gain in the realistic system.

As an example we consider Er^{3+}:YAG ($n_{\rm bg} = 1.82$) where the $4I_{9/2}$ to $4I_{15/2}$ ($\gamma_{21}^{\rm rad} = 45$ Hz) transition at 813.2 nm (transition 2-1 in Fig. 4) has a closely matched excited state absorption transition (transition 2-3 in Fig. 4) from $4I_{9/2}$ to $4G_{9/2}$ ($\gamma_{32}^{\rm rad} = 15$ Hz) with $\omega_{32} - \omega_{21} =$ 20 GHz [13,14]. Coherent driving of the transition between the next Stark level of the ground state and $4I_{9/2}$ level (transition 2-0 in Fig. 4) can be used for modulation of level 2 position, while coherent driving of $4I_{15/2}$ and $4G_{9/2}$ can be used for matching of the parameters of the upper and lower transitions in the effective ladder system. Taking $N = 1.4 \times 10^{21}$ cm^{-3} and low enough temperature to limit phonon broadening we assume $\gamma_{32} = 0.8$ GHz, $\gamma_{21} = 0.3$ GHz, and $\gamma_{42} = 0.2$ GHz. Choosing pump parameter $p = 0.035$ and the following parameters of the driving fields: $\Omega_s = 2.45$ GHz, $\Delta_s = 10$ GHz, $\Omega_c = 7.449$ GHz, and $\Delta_c = 17.893$ GHz, we obtain 3.3% refractive index modulation with respect to background value ($\Delta\chi' = 0.22$) with a periodically modulated practically vanishing absorption (max $|\chi''| < 0.0033$) as shown in Fig. 5. This result follows from the numerical analysis of the 5 level system driven with two coherent fields, and is well approximated by the analytical formula in Eq. (5). We note that the chosen wavelength mismatch, $\lambda_s - \lambda_p = 1.45$ nm, is much smaller than the width of the Bragg band gap, $\lambda\Delta n/(\pi n_{\rm bg})$, which in our case is equal to 8 nm. Already a relatively thin medium with $L = 100\ \mu$m (which corresponds to 245 periods of modulation) provides a quite high reflection coefficient, $R = 0.99998$. As the probe field is detuned from atomic resonance there will be absorption or gain which alternates on the scale of the wavelength as shown in Fig. 3, resulting in zero net absorption or gain. The produced DBR has a very narrow bandwidth of 0.6 GHz (defined by the linewidth of atomic resonance) and may be used as a frequency selective reflector.

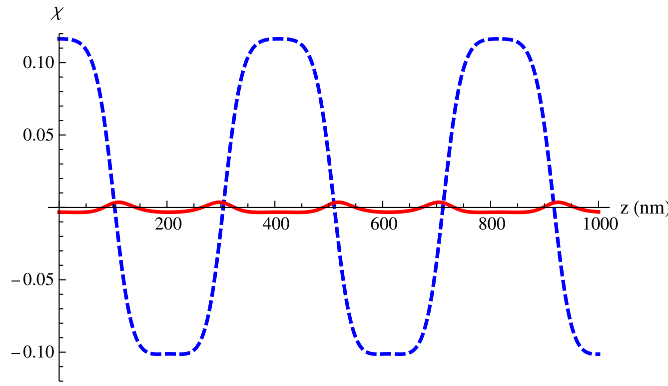

FIG. 5 (color online). Real (dashed line) and imaginary (solid line) part of the susceptibility as a function of distance along the optical axis for implementation of a optically controlled distributed Bragg grating in Er^{3+}:YAG with the parameters listed in the Letter.

In conclusion, we proposed a method to produce periodic modulation of the refractive index while keeping zero net absorption or gain. The method is based on spatial modulation of the energy of the populated intermediate state in an effective three-level system with matched transition properties by an external strong control field via the ac-Stark effect. Possible implementation of this technique in Er^{3+}:YAG is suggested, where a 3% modulation of the refractive index with vanishing absorption is possible. The proposed method may find useful applications for the creation of optically controllable photonic structures such as distributed Bragg reflectors, holey fibers, photonic crystals, etc. A major advantage of these structures as compared to traditional photonic structures is that they can be easily manipulated (including switching on or off, changing the amplitude and period of modulation) by varying the parameters of the optical control fields.

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*cobrien.physics@gmail.com

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