where the field is approximated by a Gaussian function, the parameters are homogeneous in all the coordinates. The total strain is given by equation (1) and (2) where the strain due to the applied force is given by equation (3).

\[
\varepsilon_c = \varepsilon_0 + \varepsilon_1 + \varepsilon_2 + \varepsilon_3 + \varepsilon_4 + \varepsilon_5
\]

The total strain is given by the sum of the applied force and the internal stresses. The internal stresses are due to the interaction between the lattice and the applied force. The applied force is given by the equation (4) and the internal stresses are given by equation (5) and (6).
\[ \begin{align*}
E_3 &= \exp \left( -\frac{\alpha_2}{2} \right) \left[ \gamma_{23} R E_{23} \sinh(Rz) \right], \\
E_4 &= \exp \left( -\frac{\alpha_2}{2} \right) \left[ \gamma_{24} R E_{23} \sinh(Rz) \right].
\end{align*} \]

Here \( R = \sqrt{\beta^2 + \gamma^2}, \beta = \frac{(\alpha_4 - \alpha_2)(i/2 + iA\Delta k)}{2}, \Delta k = \delta k_1 + \delta k_2 - \delta k_3 - \delta k_4, \gamma = \gamma_{24} \gamma_{43}, \gamma_{42} = \gamma_{43} E_1 E_2, \) and \( E_{23} \) and \( E_{43} \) are input values (at \( z = 0 \)). The first terms in brackets indicate FWM and second-order optical parametric amplification (OPA) processes. If either of the driving fields is switched off, \( \gamma_{24} = 0 \), and the weak radiations are described as \( |I_{42}|^2 = |I_{40,23}|^2 \exp(-\alpha_4 L) \) \((I_{42} = |E_{42}|^2)\). Owing to FWM coupling, probe radiation \( E_4 \) generates a wave \( E_2 \) close to \( \omega_{2p} \), which in turn contributes to \( E_4 \) because of FWM. This process results in correlated propagation of two waves along the medium. A gain or absorption of any of them influences the propagation features of the other. If an absorption (gain) exceeds the rate of the FWM conversion \((|\gamma_4|^2/|\beta^2| < 1)\), \( \Delta k = 0 \); if \( E_{23} = 0 \), and \( E_{23} \neq 0 \), we obtain at \( z = L \) the result \( I_2/I_{23} = \exp(-\alpha_4 L/2) + |\gamma_2|^2/(\beta^2)^2 \exp(-\alpha_2 L/2)|^2 \), where \( \alpha_2 = -\alpha_4 \). Alternatively, if \( E_{23} = 0 \), and \( E_{23} \neq 0 \), \( I_2/I_{23} = |\gamma_2|^2/(\beta^2)^2 \exp(-\alpha_2 L/2)|^2 \). One can see that achieving gain requires large optical lengths \( L \) and significant Stokes gain on the transition \( \gamma_2 \) \((|\gamma_2 L/2| \gg |\beta|^2 \gamma_2^2)\), as well as effective FWM at both \( \omega_2 \) and \( \omega_4 \). The dependence \( I_2(L) \) is predetermined by the sign of \( \Im \gamma_{24} \) and \( \Re \gamma_{42} \).

The important feature of the far-from-degenerate interaction is that the magnitude and the sign of the multiphoton resonance detunings and, consequently, of the amplitude and phase of the lower-state coherence \( \rho_{21} \) differ for molecules at different velocities because of the Doppler shifts. Such is not the case in near-degenerate schemes. The interference of elementary pathways, with Maxwell’s velocity distribution and saturation effects taken into account, results in a nontrivial dependence of the microscopic parameters on the intensities of the driving fields and on the frequency detunings from the resonances \([9]\). A density matrix solution was found exactly with respect to \( E_{13} \) and a first approximation for \( E_{23} \). We use these formulas here for numerical averaging over velocities, for analysis of the quantities \( \alpha_4, \gamma_{42} \), and also to obtain a numerical solution of the system of Eqs. \((1)-(2)\), with the inhomogeneity of the coefficients taken into account. We stress that the effect under consideration is AVI rather than conventional OPA accompanied by absorption and Stokes gain, because the quantum interference involved in resonant schemes is so crucial a role that thinking in terms of the Manley-Rowe conservation law would be misleading \([10]\).

The main outcomes of the simulations the conditions of the experiment are illustrated in Figs. 1(b), 2, and 3. The transitions of Fig. 1(a) and relaxation parameters are attributed to those of N42; \( \lambda_{41-4} = 655, 756, 532, 480 \) nm, \( \Gamma_m, g, n = 260, 200, 30, \gamma_{mn, ml, gn, gt} = 24, 20, 10, 40, \Gamma_{mn, in, on, gn, gt} = 110, \Gamma_{mn} = 130, \) and \( \gamma_{ng, ij} = 140 \) (all in \( 10^3 \times s^{-1} \)). Here \( \Gamma_i \) is the population, \( \Gamma_{ij} \) is the coherence, and \( \gamma_{ij} \) are the spontaneous interlevel relaxation rates. At \( T = 450^\circ C \), the Doppler FWHM of the transition at \( \lambda_4 \) is 1.7 GHz, and the Boltzmann population of level \( n \) is 2% of that of level \( l \). The Rabi frequencies \( G_1 = \Gamma_1 \delta_1/2 \gamma \) and \( G_3 = \Gamma_3 \delta_3/2 \gamma \) of \( \sim 100 \) MHz correspond to 100-mW beams focused on a spot with sizes of a few parts of a millimeter, i.e., one photon per several molecules. However, the presence of such fields give

\[ \begin{align*}
E_2 &= \exp \left( -\frac{\alpha_2}{2} \right) \left[ \gamma_{23} R E_{23} \sinh(Rz) \right], \\
E_4 &= \exp \left( -\frac{\alpha_2}{2} \right) \left[ \gamma_{24} R E_{23} \sinh(Rz) \right].
\end{align*} \]
imaginary parts take on even different signs [as in Fig. 2(c)-2(f)]. This behavior is in marked contrast to that of solid-state and off-resonant nonlinear optics.

The inhomogeneity of the driving fields [Fig. 3(a)] gives rise to a significant change of the material parameters along the medium (dashed curves in Fig. 2), so $\alpha_3$ may even increase above its value in the weak-field limit. The interplay of these effects determines the spatial dynamics, optimum parameters, and achievable gain [Fig. 1(b)]. Along a substantial medium length, the probe field is only being depleted [Fig. 3(b)]. Its growth begins at the length where the generated and enhanced field $E_2$ [dashed curves in Fig. 3(b)] becomes comparable with $E_{48}$. The simulations explicitly reveal that the fully resonant conditions explored in Ref. 8 are far from optimal [inset in Fig. 3(b)], and most probably the gain reported in Ref. 8 is a misinterpretation of the experiment. The maximum gain in Fig. 1(b) is 1050, which is well above the characteristic threshold for self-oscillation to be established inside the optical cavity from the spontaneous radiation. This gain can readily be increased further to the mirrorless oscillation level. Both linear and laser-induced nonlinear dispersion inhomogeneous along the medium are taken into account in Figs. 1(b) and 3. Our results also demonstrate that the problem of AWI in similar schemes may not be reduced to the condition of a sign change of $\alpha_3$, as was done in the research reported in Ref. 11. The solid curve in Fig. 3(b) shows that there is an optical thickness controlled by the driving radiations whose small variation results in a switching from the absorption regime to transparency and further to amplification. Figure 3(c) presents the possibility of controlling this switching with a small change of either the frequency of the probe radiation or the intensity of the driving radiation (inset, Fig. 3). Obviously, the same processes can be employed for generating and manipulating large dispersion without the accompanying depletion of radiation. The required intensity can be further decreased in identical but more favorable atomic schemes. Owing to the generated molecular coherence, fields $E_3$ and $E_2$ may possess nearly perfect quantum correlations that yield almost complete squeezing[6].

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