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Quasisimultons in Thermal Atomic Vapors

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The propagation of two-color laser fields through optically thick atomic ensembles is studied. We demonstrate how the interaction between these two fields spawns the formation of copropagating, two-color solitonlike pulses akin to the simultons found by Konopnicki and Eberly [Phys. Rev. A **24**, 2567 (1981)]. For the particular case of thermal Rb atoms exposed to a combination of a weak cw laser field resonant on the $D1$ transition and a strong sub-ns laser pulse resonant on the $D2$ transition, simulton formation is initiated by an interplay between the $5s_{1/2} - 5p_{1/2}$ and $5s_{1/2} - 5p_{3/2}$ coherences. The interplay amplifies the $D1$ field at the arrival of the $D2$ pulse, producing a sech-squared pulse with a length of less than 10 μm. This amplification is demonstrated in a time-resolved measurement of the light transmitted through a thin thermal cell. We find good agreement between experiment and a model that includes the hyperfine structure of the relevant levels. With the addition of Rydberg dressing, quasisimultons may offer interesting prospects for strong photon-photon interactions in a robust environment.

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Self-induced transparency manifests dramatically by the formation of optical solitons propagating undistorted over long distances in a medium opaque to a cw field of the same wavelength [1]. A short light pulse may propagate as a soliton or split into multiple solitons only if it is sufficiently intense [2]. However, it has been known for some time that a weak solitonlike pulse at one frequency can copropagate with a stronger solitonlike pulse at another frequency. Solutions of the Maxwell-Bloch equations describing this situation in doubly resonant three-level V systems were first found in the form of matched sech pulses, or simultons, under the condition that the two transitions have the same oscillator strength [3,4]. Remarkably, the two pulses may copropagate as a simulton even if neither of them is strong enough to support a soliton in the absence of the other field [5]. The condition that the two transitions have the same oscillator strength is difficult to realize experimentally [6] but makes the Maxwell-Bloch equations integrable in the sense of the inverse scattering transform (in a suitable approximation), which permits an in-depth analysis of their analytical solutions [7].

However, this condition can be relaxed without compromising the formation of pairs of solitonlike pulses copropagating with little distortion over much longer distances than allowed by Beer's law [8,9]. We refer to such pairs of pulses as quasisimultons, to distinguish them from the ideal sech simultons of Ref. [3]. It has been noted, in particular, that a soliton on one transition of a V system may enhance transmission of a weak pulse on the other transition even if the latter has a different oscillator strength [10]. The soliton may amplify the weak pulse and transport it through the medium simultonlike [11,12]. The term soliton-induced transparency has been coined for this effect [11]. To our knowledge, it has previously been observed only in the propagation of super-radiance pulses in a neon plasma [13].

In this Letter, we show that soliton-induced transparency is readily seen in experiments using thermal vapor cells, in our case a thin cell containing a rubidium vapor [14,15] addressed by a cw field resonant on the $D1$ transition and a pulsed field resonant on the $D2$ transition. The numerical simulations described below reproduce the observed increase in transmission and pulse shaping indicating that this change signals the formation of a quasisimulton.

A sketch of the experimental scheme is shown in Fig. 1(a). A dense, 2-μm thick thermal vapor of rubidium atoms in their natural isotopic abundances interacts with two copropagating monochromatic laser beams forming a V-type excitation scheme. The probe and coupling beams are linearly polarized in orthogonal directions. They are resonant on, respectively, the $5s_{1/2}(F=3)-5p_{1/2}(F=3)$ and $5s_{1/2}(F=3)-5p_{3/2}(F=4)$ transitions of ^{85}Rb. The coupling beam is focused to a waist of ~20 μm while the probe beam is focused more tightly to a waist of ~10 μm, which minimizes variation of the coupling intensity for the atoms in the probe beam. The probe field applied to the cell is cw. The coupling field is shaped to a short, nearly Gaussian pulse of a duration of typically 0.8 ns full width at half maximum (FWHM). Taking losses into account, we estimate that at the front of the medium, and on axis, the coupling field had an intensity of 3.7 kW cm^{-2} at a measured peak power of 85 mW and the probe field an intensity of 24 W cm^{-2}. Following propagation the two

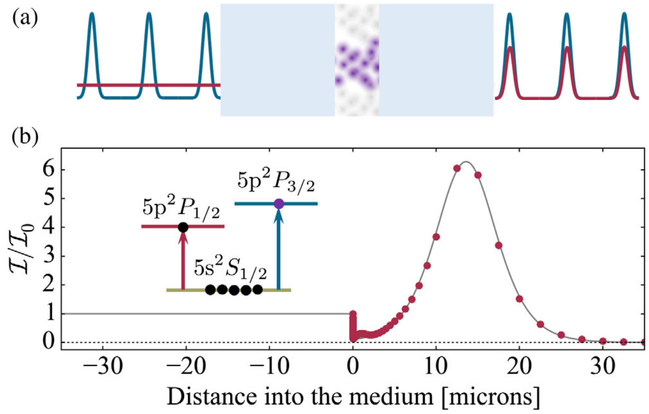

FIG. 1. (a) Schematic of the experiment (not to scale). A cw probe field (red) resonant with the $5s_{1/2}(F=3)-5p_{1/2}(F=3)$ transition of ^{85}Rb copropagates with a pulsed coupling field (blue) resonant with the $5s_{1/2}(F=3)-5p_{3/2}(F=4)$ transition. The coupling field reshapes the probe into sech-squared pulses. (b) Inset: The level scheme indicating the probe and coupling fields. Main plot: A snapshot of the intensity of the probe field (relative to the incident intensity) calculated over a much longer propagation distance than achieved in the experiment, taken 0.05 ns after the coupling pulse reached its maximum at the front of the medium; the circles are calculated points and the solid line is a fit of these with a sech-squared pulse of width 8.6 μm (FWHM).

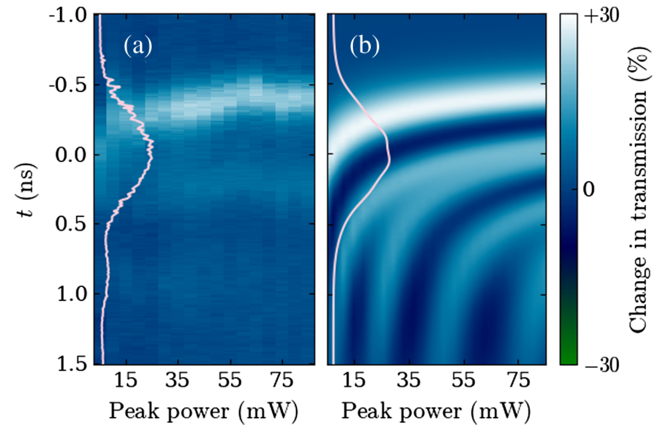

FIG. 2. The change in the transmission at the probe frequency for a 2 μm-long Rb vapor relative to the average transmission at $t<-1$ ns, vs the peak power of the incident coupling pulse. The pink curves represent the temporal intensity profile of the latter, at the back of the medium, for an initial peak power of 85 mW. The 0 of the color scale corresponds to no change compared to the transmission before the arrival of the coupling pulse. The temperature is 200 °C. (a) As measured. (b) As predicted by the model described in the text.

fields are separated by a polarizing beam splitter and their transmission through the medium is monitored using a fast photodiode [16,22]. The temporal variation of the measured probe field for various peak powers of the coupling pulse is shown in Fig. 2(a). The main feature of these data is a strong increase in transmission on the raising edge of the coupling pulse.

The other panel of this figure shows the results predicted by a model described below. Before addressing these, we first consider a simplified model consisting of a single ground state (state 0) resonantly coupled to a first excited state by a weak field (the probe field) and to a second excited state by a strong pulse (the coupling field) [Fig. 1(b)]. We take the coupling field at the front of the medium to be a Gaussian pulse of 1 kW cm^{-2} peak intensity and 0.8 ns duration (FWHM in intensity), and the probe field to have an initial constant intensity of 10 μW cm^{-2}. We set the excited state lifetimes and the transition wavelengths and dipole moments to values corresponding to the $D1$ and $D2$ lines of ^{85}Rb, respectively. The oscillator strengths of the two transitions thus differ by a factor of 2 here. We neglect Doppler broadening, self-broadening, and the hyperfine structure of the states for the time being. We assume one-dimensional propagation [23], make the rotating wave and slowly varying envelope approximations, and solve the resulting Maxwell-Bloch equations numerically, taking the atoms to be initially in the steady state driven by the probe field.

Figures 1(b) and 3 show how the probe and coupling fields vary both in space and in time within this three-state model. Figures 3(a) and 3(b) refer to the probe field. As seen from these figures, this field practically vanishes as soon as it enters the medium, prior to the arrival of the coupling pulse (the attenuation is extremely fast because the field is resonant with the transition and Doppler broadening is neglected). However, the arrival of that pulse triggers a more complicated dynamics. Microscopically, the onset of this dynamics can be traced to a rapid increase of the ρ_{12} coherence. This increase, in turn, produces a large variation and a change of sign of the ρ_{01} coherence, leading to an amplification of the probe field without a population inversion between the ground state and the first excited state [10,13]. In particular, the probe field develops into three successive pulses penetrating far deeper into the medium than the initial cw field. As shown in Fig. 3(a), these three pulses propagate solitonlike over many tens of micrometers, each at a different speed (the second of these three pulses is almost invisible in the figure). The first one is the fastest, although its maximum speed is only about $2\times10^{-4}c$. It is also the strongest, and at its peak reaches an intensity larger than that of the incident cw field by almost a factor of 7. Like the other two, it becomes weaker and slower as it propagates. A snapshot of the spatial profile of the probe field at a time when this first pulse has just formed is shown in Fig. 1(b). The strong spatial localization of this pulse caused by the slow light effect is worth noting, as is its almost pure sech^2 profile.

The 0.8 ns coupling pulse also splits into three solitonlike pulses; the first and second ones are well visible in Fig. 3(c), but not the third one. This pulse is strong enough

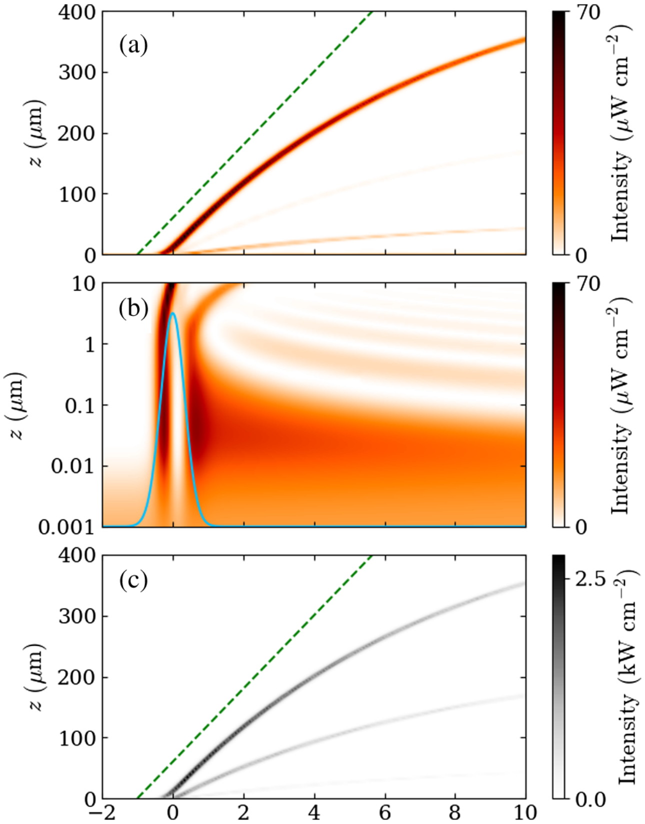

FIG. 3. Formation of quasisimultons in a three-state model of an isotopically pure ^{85}Rb vapor, neglecting Doppler broadening and self-broadening. The atomic density corresponds to a temperature of 220 °C. (a) The intensity of the probe field vs time and propagation distance within the medium. This field is resonant on the $D1$ transition and has a constant intensity of 10 μW cm^{-2} at the entrance of the medium ($z = 0$). The 0.8-ns coupling pulse is resonant on the $D2$ transition and has a peak intensity of 1 kW cm^{-2} at $(t = 0, z = 0)$. The dashed green line indicates the trajectory that a pulse propagating at a constant speed of $2 \times 10^{-4}c$ would trace in the figure. (b) Enlargement of the small-z region of panel (a). The blue curve represents the temporal profile of the coupling pulse (the intensity scale is arbitrary). (c) The same as (a) but for the coupling field.

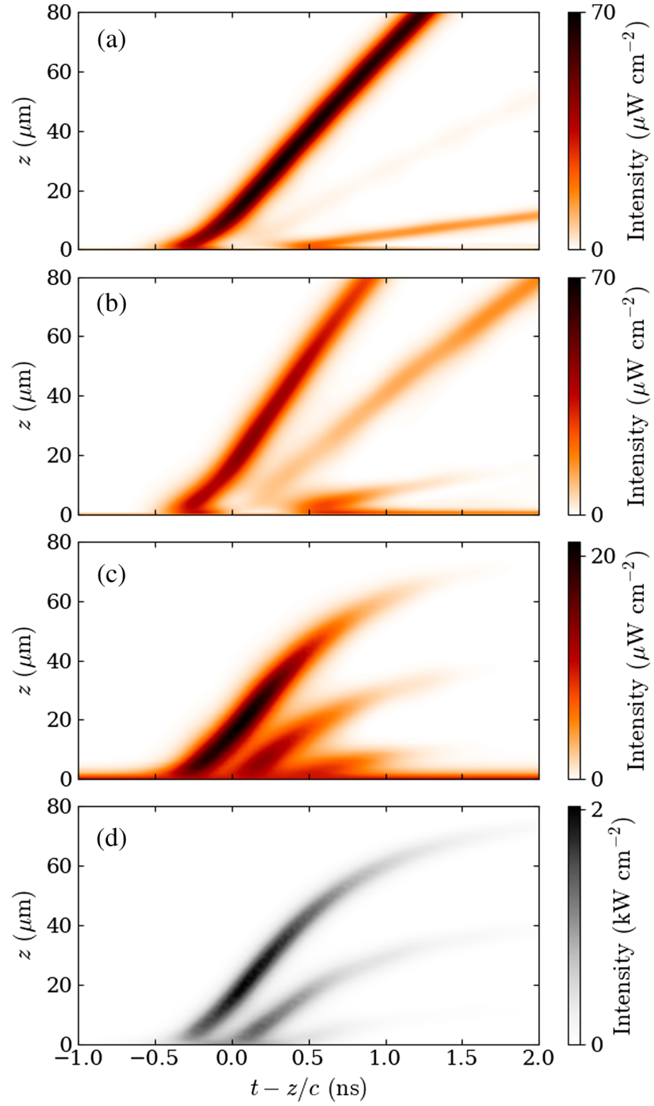

FIG. 4. (a) The intensity of the probe field in the three-state model of Fig. 3. This part of the figure is an enlargement of the small $|t - z/c|$, small-z region of Fig. 3(b). (b) The same as (a) but with the full hyperfine structure of the $5s_{1/2}$, $5p_{1/2}$, and $5p_{3/2}$ states included in the calculation. (c) The same as (b) but now also including Doppler broadening and self-broadening. (d) The same as (c) but for the coupling field.

to break into three solitons even in the absence of the probe field. The presence of the latter does not affect the propagation of the coupling field significantly, but each of these three solitons copropagates with a pulse at the probe frequency, thereby forming three quasisimultons.

Differences in transition dipole moments between the magnetic substates coupled to each other by the two fields may compromise the formation of simultons when the relevant levels have a hyperfine structure. However, this issue is not of major importance here because the bandwidth of the coupling pulse is sufficiently large compared to the energy splitting between the relevant hyperfine states [24]. Figure 4 compares the predictions of the three-state model, in (a), to the results obtained when the complete hyperfine structure of the relevant levels is included in the calculation, in (b). While there are differences between these two sets of results, it is clear that in the present case the hyperfine structure does not prevent the formation of quasisimultons and their propagation over a considerable distance. However, the hyperfine structure of the $5s$ and $5p$ levels is an important issue for longer coupling pulses [16,25].

We now include not only the hyperfine structure but also Doppler broadening and self-broadening. The corresponding results are shown in panels (c) and (d) of Fig. 4. Quasisimultons are still found in these more complete calculations. Although not as stable, they still propagate over far longer distances than would be the case for weak cw fields of the same wavelengths, and the probe field is still amplified through its interaction with the coupling pulse.

We used the same model as in Figs. 4(c) and 4(d) to obtain the results displayed in Fig. 2(b), except that we assumed the incident probe field had the same (much higher) intensity as in the experiment. We ignored the transverse Gaussian profiles of the laser beam, as factoring these in would have led to excessively long computations. Interaction with the windows was taken into account by broadening each state by 30 MHz [26]. Comparing Fig. 2(b) to Fig. 2(a), we see that the model does not predict the rapid damping of the dynamics that follows the initial increase in transmission on the raising edge of the coupling pulse (the origin of this damping is as yet unknown). However, it reproduces this strong increase well.

We also ran this calculation for a 50-μm-long cell, so as to see how this increase in the transmission develops over longer propagation distances (Fig. 5). Taking pulse reshaping into account would be necessary for comparing to measurements for a cell of that length; however, here we only aim at illustrating how the one-dimensional dynamics would evolve beyond 2 μm if it remained unperturbed. As seen from the figure, this enhancement would develop into a well-defined pulse copropagating with the first of the solitons the coupling pulse splits into. It can thus be identified with the formation of a quasisimulton.

Although the probe field is considerably more intense in Fig. 5 than in Figs. 3 and 4, the pulses at the probe frequency still closely follow the pulses at the coupling frequency and the latter still propagate essentially as if the probe field was absent. As an aside, we note that the presence of the probe field starts affecting the coupling field at still higher probe powers than considered here. The ensuing dynamics will be discussed elsewhere.

In conclusion, we have demonstrated the amplification of a weak probe field by a strong pulse—the first step in the formation of quasisimultons. Such simultons allow the propagation of even arbitrarily weak localized fields through optically thick media (at least for classical fields). The satisfactory agreement between theory and experiment found in Fig. 2 shows the applicability of our model, extended as necessary, to the exploration of what other two-color propagation phenomena could be observed in thermal atomic vapors. Future work will focus on photon-photon interactions. Because of the slow light effect (propagation speeds of order $c/5000$) and strong focusing (a transverse beam size of order 2.5 μm gives a Rayleigh range over two times the length of the simulton), a high optical Kerr nonlinearity is expected. In addition, as the simulton length (<10 μm) is close to the Rydberg blockade radius [27], it is interesting to investigate simulton-simulton collisions in the presence of Rydberg dressing. As in the context of Rydberg quantum optics [28], there is the potential that the nonlinearity will be significant at the few photon level.

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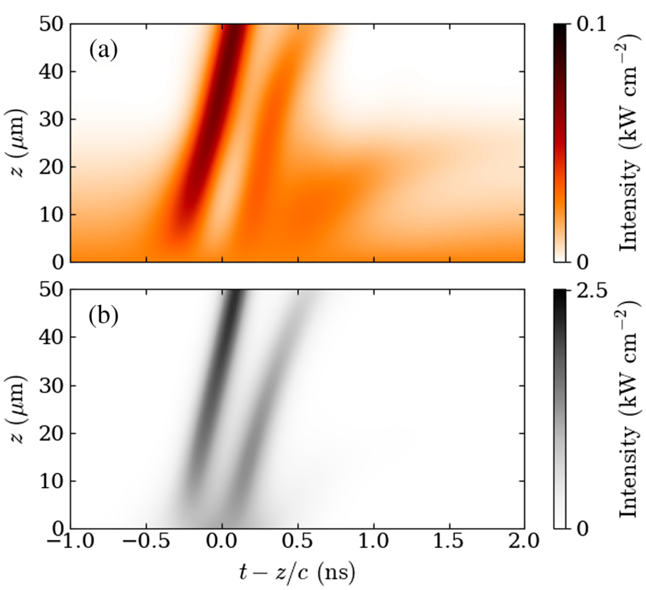

FIG. 5. (a) The calculated probe field for the conditions of Fig. 2, now for a much longer propagation distance of 50 μm. The peak power of the coupling pulse is 35 mW. (b) The corresponding coupling field.

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$$\theta = \frac{2d}{\hbar}\int_{-\infty}^{\infty}\mathcal{E}(t)dt.$$

A pulse reshapes into a superposition of N 2π-solitons when initially $(2N-1)\pi < \theta < (2N+1)\pi$ [1].

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