Controlling Light by Light with an Optical Event Horizon

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A novel concept for an all-optical transistor is proposed and verified numerically. This concept relies on cross-phase modulation between a signal and a control pulse. Other than previous approaches, the interaction length is extended by temporally locking control and the signal pulse in an optical event horizon, enabling continuous modification of the central wavelength, energy, and duration of a signal pulse by an up to sevenfold weaker control pulse. Moreover, if the signal pulse is a soliton it may maintain its solitonic properties during the switching process. The proposed all-optical switching concept fulfills all criteria for a useful optical transistor in [Nat. Photon. **4**, 3 (2010)], in particular, fan-out and cascadability, which have previously proven as the most difficult to meet.

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Manipulating light with light by interaction of optical pulses in a nonlinear medium is an active area of research, in particular with respect to optical transistors [1]. In an optical fiber with cubic nonlinearity, two copropagating pulses interact via cross-phase modulation (XPM) [2]. Many different approaches exist to exploit this effect for an all-optical control of optical light pulses; see, e.g., [3] where it has been shown how a weak signal pulse (SP) can effectively be compressed by a strong control pulse (CP). All-optical switching based on XPM has also been discussed in [4,5], however, with the same requirement of a strong CP. Moreover, no optical transistor has been demonstrated to date that can be cascaded in several stages. As a consequence of these and other shortcomings, Miller defined a set of arduous criteria that he considered mandatory for a practical all-optical switching device [1], with no currently known optical transistor concept complying to all criteria simultaneously.

In response to this seemingly unsolvable quest, we exploit an optical event horizon for extending the effective interaction length of two pulses copropagating in a nonlinear optical fiber. This recently discovered phenomenon [6] is an unexpected analogy between nonlinear fiber optics and trapping of light in gravitational fields. In the context of supercontinuum generation [7], this mechanism may lead to energy localization or trapping in the normal dispersion regime [8]. With an appropriate choice of the initial conditions, this gravitylike effect has been used in [9] to demonstrate propagation dynamics similar to that in quantum bouncing [10]. In the optical analogue, the bouncing affects reflection of a dispersive wave packet off an accelerating soliton. Concomitant with the soliton being bounced, the frequency of the radiation is up-shifted. This reflection of radiation and the associated blueshift have also been observed in [6] and were interpreted as the fiber-optical analogue of a white-hole event horizon, leading to a conceptually similar frequency shift as in astrophysics. Using an optical setup, one can directly observe Hawking's radiation at the phase velocity event horizon [11]. The latter should be distinguished from the group-velocity event horizon (see [12] and below).

Here we show that reflection of a dispersive wave packet at the group-velocity horizon of a fundamental soliton represents a strong light-light interaction in an optical fiber. Combined with a strongly curved anomalous dispersion profile, it can be exploited for efficient control of an optical pulse by another optical pulse. We modify the duration, intensity, and carrier frequency of a SP by suitably choosing parameters of the CP in the normal dispersion regime. In particular, our concept fulfills all criteria for a useful all-optical transistor.

In the following, we consider the case that the CP slowly passes the SP with carrier frequencies ω_c and ω_s, respectively. Ideally, the condition $n_g(\omega_c) + \delta n > n_g(\omega_s) > n_g(\omega_c)$ upon the optical group-velocity horizon should be fulfilled, where $n_g(\omega)$ is the group index and δn is the nonlinear refractive change due to the strong SP. Initially, for vanishing temporal overlap and no resulting XPM, the CP approaches the SP from positive delays. When the pulses begin to overlap, an extended XPM interaction builds up, preventing the pulses from crossing each other, and the approach is eventually halted. The two pulses may be temporally locked due to their mutual interaction. Such extended interaction considerably changes properties of both pulses. An effective interaction is achieved even when a weaker condition is fulfilled, $n_g(\omega_s) \gtrsim n_g(\omega_c)$, such that the velocity of the CP is somewhat larger than that of the SP. A more interesting scenario arises when the

SP slowly passes the CP (see below). In both cases, the interacting optical fields contain considerably different frequencies ω_c and ω_s, such that the slowly varying envelope approximation (SVEA) may be violated. For simplicity, we consider a single-mode waveguide. The optical field is characterized by a single real-valued component $E(z, t)$ which we write as a discrete sum in the spectral domain $E(z, t) = \sum_\omega E_\omega(z)e^{-i\omega t}$. Here $\omega T \in 2\pi\mathbb{Z}$, T is a large period in the time domain, z refers to the propagation distance, and any dependence perpendicular to z is integrated out. Addressing ultrashort pulse propagation, we introduce a complex-valued analytic signal $\mathcal{E}(z, t) = 2\sum_{\omega>0}E_\omega(z)e^{-i\omega t}$, consisting of summation of spectral components with positive frequencies. $\mathcal{E}(z, t)$ is subject to the following model equation [13]:

$$i\partial_z\mathcal{E}_\omega + \beta(\omega)\mathcal{E}_\omega + \frac{3\omega^2\chi^{(3)}}{8c^2\beta(\omega)}(|\mathcal{E}|^2\mathcal{E})_{\omega>0} = 0, \quad (1)$$

where $\chi^{(3)}$ refers to the Kerr nonlinearity. If SVEA with respect to the SP applies, one can reduce (1) to the nonlinear Schrödinger equation [2] with the nonlinearity parameter $\gamma = (3\omega_s\chi^{(3)})/[4\epsilon_0c^2n^2(\omega_s)A_{\rm eff}]$, where $A_{\rm eff}$ is the effective fiber area.

In our investigations, we intentionally exclude all terms that may hide the main ingredients in the scattering process at the optical event horizon. In particular, we exclude Raman scattering, which induces a soliton self-frequency shift, leading to a strong influence on its group velocity. Equation (1) directly delivers the electric field $E = {\rm Re}\mathcal{E}$ and avoids use of the envelope. We also use an improved description of the propagation constant $\beta(\omega)$. Here, the common polynomial expansion involves too many terms and may cause numerical stiffness. A rational approximation is used instead that leads to a better approximation of $\beta(\omega)$ for higher frequencies [14].

Our main precondition is the copropagation of a CP (normal dispersion regime) and a fundamental signal soliton (anomalous dispersion) at nearly identical group velocity. Consequently, the effective XPM interaction length between the two pulses significantly increases and walk-off effects diminish. For a manipulation of an intense fundamental soliton by a weak CP, one needs a further important condition. The dispersion in the vicinity of the central frequency of the soliton has to vary significantly. This condition is fulfilled by particular frequency combinations, given an adequate refractive index profile, as demonstrated by the example of fluoride glass. The according group-velocity dispersion β_2 and the relative group delay $\beta_1 = 1/v_g$ are shown in Fig. 1. To trap a weaker CP at the optical event horizon, the fundamental soliton has to exhibit a fast transient intensity, which is assured by ultrashort pulse durations in the sub-100 fs regime. We launch two pulses,

$$\mathcal{E}(z = 0, t) = \frac{\mathcal{A}_c e^{-i\omega_c t}}{\cosh[(t+\delta)/t_c]} + \frac{\mathcal{A}_s e^{-i\omega_s t}}{\cosh[(t-\delta)/t_s]},$$

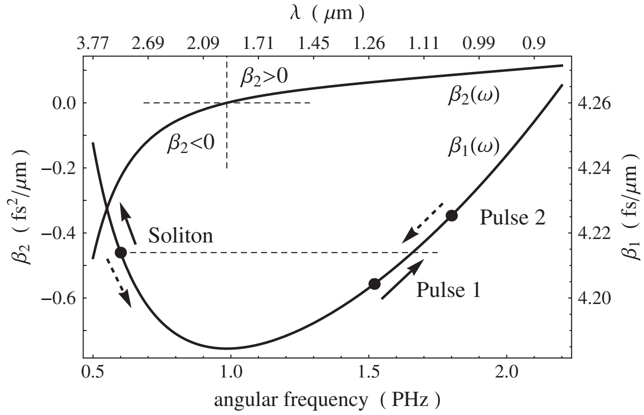

FIG. 1. Group-velocity dispersion β_2 and relative group delay β_1 of fluoride glass, exhibiting frequency combinations for a soliton and a CP with equal group velocities. The soliton is injected at $\omega_s = 0.6$ PHz. Arrows indicate the induced frequency shift after the collision with a faster and a slower CP at $\omega_c = 1.52$ PHz (solid arrow) and 1.8 PHz (dotted arrow), respectively.

into the fiber and, on either side of the zero dispersion wavelength, we synchronize their time delays 2δ at nearly equal group velocities to meet the above-discussed conditions. The launched pulse spectra are strongly separated, ensuring vanishing overlap between them within orders of magnitude. The amplitude $\mathcal{A}_c$ of the CP is significantly lower than that of the signal soliton $\mathcal{A}_s$, such that we only have to consider the intensity dependent refractive index change induced by the soliton. For the observation of a wavelength shift towards shorter ("bluer") wavelengths of the dispersive radiation [6,9], we inject the soliton prior to the CP into the fiber with a delay of 400 fs. For the present dispersion profile, $\omega_s = 0.6$ PHz and $\omega_c = 1.52$ PHz are a suitable choice for the soliton and CP, respectively. The width t_s of the former at $\beta_2(\omega_s) = -0.229$ fs$^2/\mu$m has to be chosen such that the intensity induces effectively a small yet sufficient increase of the refractive index to build up a fiber-optical group-velocity event horizon, involving a significant change of group velocities. We start with an input soliton pulse width of $t_s = 21$ fs, a pulse width of the dispersive wave $t_c = 70$ fs, and a 9:1 peak power relation between these two pulses. For a nonlinear fiber with $\gamma = 0.1$ W^{-1} m^{-1} this corresponds to peak powers of $P_s = 5.2$ kW and $P_c = 0.58$ kW for the soliton and CP, respectively.

The temporal propagation dynamics are shown in Fig. 2(a) in a reference frame moving with an unperturbed soliton at ω_s. Because of group-velocity dispersion, the CP approaches the soliton, until it reaches the trailing edge of the soliton. When the pulses begin to interact, XPM induces a substantial frequency shift of their central frequencies [15]. The central frequencies of the trailing and leading pulse shift to the blue and the red, respectively [Fig. 2(b)]. However, contrary to the typical XPM action

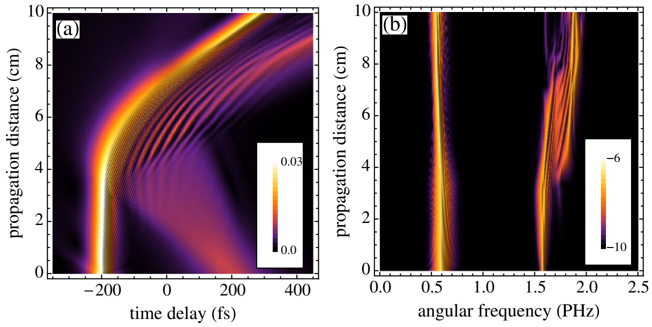

FIG. 2 (color online). Time domain (a) and spectral (b) evolution along the fiber with a signal at $\omega_s = 0.6$ PHz and a faster CP 1 at $\omega_c = 1.52$ PHz, representing a typical scattering process of a dispersive wave at an optical horizon of a soliton. Energy is transferred from the soliton to the dispersive wave.

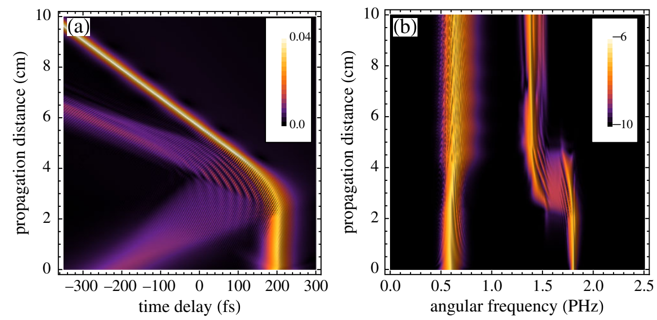

FIG. 3 (color online). Time domain (a) and spectral (b) evolution along the fiber with a signal at $\omega_s = 0.6$ PHz and a slower CP 2 at $\omega_c = 1.8$ PHz, representing a scattering process with a transfer of energy from the dispersive wave to the soliton.

the CP cannot pass the soliton. Because of strong nonlinear refractive index changes, the weak pulse is reflected at the trailing edge of the soliton, and the nonlinear interaction persists only between the weak pulse and the trailing edge of the strong pulse. After reflection, both pulses thus drift apart with smaller group velocities than before. An important feature here is that the soliton SP is strongly affected by the reflection process. Figure 2(b) shows the CP frequency shift towards shorter wavelength and an adverse shift towards longer wavelengths of the soliton pulse. In [6,9] the redshift of the soliton center wavelength and the change of the soliton group velocity interfere with the Raman effect. In the absence of the Raman effect, Fig. 2(b) can then be understood as an energy transfer between the pulses during collision. Given energy conservation, the weaker pulse experiences a stronger shift than the main pulse, and the latter is still shifted by more than its half-width maintaining stable propagation as a fundamental soliton. This finding opens the perspective of a solitonic transistor, enabling switching of a strong soliton pulse with a much weaker CP, as is shown in Fig. 3 and will be discussed subsequently. Now, the fundamental soliton is kept as before, but is injected after a CP with $\omega_c = 1.8$ PHz, i.e., at a slightly smaller group velocity than the soliton. The collision increases the group velocity of both pulses [Fig. 3(a)]. The central frequency of the CP is shifted to the red from $\omega_c = 1.8$ to 1.45 PHz [Fig. 3(b)]. The soliton central frequency shifts toward the blue from $\omega_s = 0.6$ to 0.62 PHz such that $\beta_2 = -0.229\ \mathrm{fs}^2/\mu\mathrm{m}$ is reduced to $-0.199\ \mathrm{fs}^2/\mu\mathrm{m}$. This leads to a soliton with double the initial peak power and a smaller duration of 11 fs.

The reflection process and the accompanied energy transfer between the pulses do not suffice to cause such a strong change in the temporal properties of the soliton as the induced shift of the soliton central frequency is small and the strong temporal change depends on the dispersion variation in the range of the induced frequency shift. This mechanism is comparable to soliton compression in a dispersion decreasing fiber, where the soliton adiabatically changes its shape due to the change of β_2 [2]. Adjusting the properties of the CP we can decrease or increase the frequency shift acting on the soliton and thereby manipulate the compression ratio. Figure 4 shows the increase of peak power and the decrease of soliton duration behind the fiber. The very mechanism behind this switching phenomenon can yet again be understood from energy conservation. An increase of intensity or pulse width of the CP leads to a stronger frequency shift for the control pulse as well as for the soliton. Consequently, the soliton frequency is shifted to a lower value of β_2, corresponding to a fundamental soliton with higher peak power. Thus the manipulation of the soliton strongly depends on the change of dispersion resulting from the frequency shift. The stronger the difference of the dispersion values, the smaller the frequency shift has to be, and vice versa. The simulations indicates that up to 10% of the energy of the fundamental soliton can be transferred between the pulses. One can also increase the difference between the input pulse velocities, leading to a higher frequency shift, due to the Doppler effect [6]. Yet any change of the control pulse parameters that can be induced to manipulate the soliton pulse is constrained by the requirement to enable the scattering process at the optical event horizon. With increasing CP energy the efficiency of the scattering process might be reduced, such that the CP is only partially reflected at the horizon,

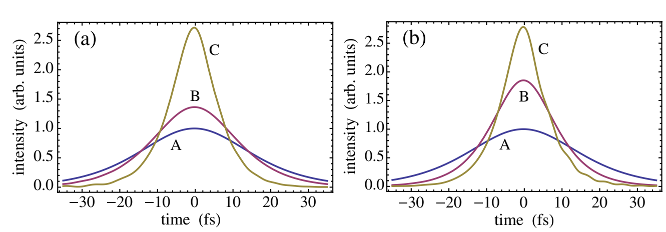

FIG. 4 (color online). Output soliton pulses (B and C shifted to $t = 0$) in relation to the input soliton pulse (A) after collision of the signal with a CP. (a) CP widths are $t_c = 70$ fs ($A \rightarrow B$) and $t_c = 90$ fs ($A \rightarrow C$). (b) Peak intensity of the soliton relates to that of the CP as 6:1 ($A \rightarrow B$) and 4:1 ($A \rightarrow C$).

with a smaller part tunneling through. Effective performance of the proposed transistor is achieved over a wide range of parameters. The reflection process represents a robust mechanism and is also observed under the impact of Raman scattering [6,9]. Considering resulting Raman deceleration effects, we observe an increase of group-velocity-matching with propagation distance, nevertheless leaving the fundamental Kerr-type scattering process mostly untouched. A detailed analysis of the influence of Raman deceleration on frequency shifts in an optical event horizon can be found in [12]. This effect was deliberately excluded in our analysis here in order to isolate the chief effect for the observed switching behavior without the necessity for dissipative mechanisms.

In conclusion, we show that the scattering of an optical pulse at an optical event horizon provides an unprecedented potential to control the properties of a light pulse with another weaker light pulse. Careful adjustment of the group velocities of the two pulses enables us to enhance their effective interaction such that their center frequencies either strongly repel or attract each other, resulting in perfectly efficient mutually induced frequency shifts. In the simplest case, the optical switching action is decoded by this frequency shift. Ensuring a suitable dispersion profile one can achieve strong changes in the soliton properties such as its width in time as well as its peak intensity. Most importantly, with proper parameter selection, a strong pulse can be switched by a 6–7 times less energetic pulse, which clearly sets our method apart from previously proposed optical transistors. Moreover, while the CP may experience serious reshaping effects, the main pulse does not dispersively spread or break up into multiple pulses, such that the solitonic switching scheme is cascadable. Our solitonic switching scheme therefore fulfills, among all others, the two most stringent criteria for an optical transistor, fan-out and cascadability [1]. Furthermore, logic-level restoration and input-output isolation can also be achieved. Operating with solitons as signals has the advantage that a nonexact fundamental soliton is changed into an exact soliton while propagating in the fiber. The degradation of the signal from an exact fundamental soliton can thereby amount to up to 50% [2]. Input and output pulses can easily be separated, as they can here easily be filtered out spectrally. In practical terms, the scheme requires the presence of a Kerr nonlinearity together with a concave group delay $\beta_1(\omega)$, cf. Fig. 1, in a reasonably long waveguide. Waveguide dispersion is helpful and required to shift the soliton frequency into the more practical telecommunications range. Apart from photonic crystal fibers, it seems appealing to investigate the use of silicon waveguides on a chip [16] for the switching, as this promises to shrink the required waveguide lengths.

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