Superradiant Amplification of an Ultrashort Laser Pulse in a Plasma by a Counterpropagating Pump

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An initially short ($<1/\omega_p$) laser pulse can be superradiantly amplified by a counterpropagating long low-intensity pump while remaining ultrashort. This superradiant amplification occurs if the frequency of the pulse is lower than that of the pump, and the initial pulse intensity is sufficiently high. Numerical simulations indicate that the short pulse can be amplified to an intensity hundreds of times the pump intensity, with the pump depletion as high as 40%. This implies that the long pump is efficiently time compressed without frequency chirping and pulse stretching, making the superradiant amplification an interesting alternative to chirped-pulse amplification. [S0031-9007(98)07762-X]

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The generation of very short ultraintense laser pulses is crucial for a number of scientific and technical applications, such as plasma-based particle accelerators, inertial confinement fusion (ICF), and x-ray lasers [1]. For example, laser-wakefield accelerators [2] require pulses of about $1/\omega_p$ duration, where $\omega_p = (4\pi e^2 n_0/m)^{1/2}$ is the plasma frequency, and n_0, $-e$, and m are the plasma density, electron charge, and mass, respectively. In one of the advanced ignition scenarios, Fast Ignitor [3], a picosecond petawatt pulse is needed to ignite the compressed fuel. The target gain improves as the energy of the ignition pulse increases [4], making a simple and efficient pulse-compression technique that is desirable for ICF applications. The standard approach to reaching extremes of power with ultrashort pulses is the chirped-pulse amplification (CPA) [5], in which a laser pulse is stretched, amplified, and recompressed. CPA-based optical systems have been shown to generate subpicosecond petawatt laser pulses [6,7] with up to 500 J per pulse. The pulse energy is limited by the thermal damage to the compression gratings which become large and expensive for kJ pulses. A possible route to increasing peak intensity is to decrease pulse duration, which can be achieved by using broad bandwidth parametric amplifiers [8]. This does not, however, ease the energy restriction.

In this Letter we suggest a novel approach to pulse compression which does not require pulse stretching or compression. We demonstrate how an ultrashort pulse can be amplified by several orders of magnitude by colliding with a long counterpropagating pumping laser in the plasma. The main advantage of using the plasma medium for amplification is that there is no thermal damage threshold—fresh plasma can be used for each shot. The counterpropagating geometry is chosen for conceptual simplicity. Different designs, such as injecting the pump from the side, can be used to reduce the level of the amplified spontaneous emission (ASE).

The approach is based on the extensively studied paradigm of three-wave interaction between the counterpropagating electromagnetic waves and a plasma wave [9–11]. The plasma wave is ponderomotively driven by the periodic intensity pattern produced by the interference between the pumping beam (PB) and the amplified beam (AB). Frequency detuning between the PB and AB determines the direction of energy flow: From Manley-Rowe AB is amplified through the Raman backscattering (RBS) instability of the pump if its frequency ω_0 is lower than the pump frequency ω_1. The interaction between the lasers is maximized when they propagate in opposite directions because the ponderomotive force which drives the plasma density perturbation is proportional to the gradient of the laser intensity. RBS instability received considerable theoretical [10–12] and experimental [13–15] attention. Earlier work on RBS focused on the amplification of long pulses growing from noise.

The novel element here is that we identified a new nonlinear regime of RBS, in which an externally injected ultrashort pulse is dramatically amplified to an intensity exceeding the pump intensity by orders of magnitude. Accessing this regime requires a sufficiently intense initial pulse: $4\omega_0^2 a_0 a_1 \geq \omega_p^2$, where $a_{0,1} = eA_{0,1}/mc^2$ are the normalized vector potentials of the AB and PB, respectively. When this condition is satisfied, electron motion is determined by the ponderomotive force and not by the space-charge electric field of the plasma wave. The resonance condition for the plasma wave excitation $\Delta\omega = -\omega_p$ does not need to be satisfied precisely.

Consider the interaction between two planar circularly polarized laser pulses $\vec{a}_0$ and $\vec{a}_1$, where $\vec{a}_{0,1} = a_{0,1}/2(\vec{e}_x \pm i\vec{e}_y)e^{i\theta_{0,1}} + \text{c.c.}$, and the phases of the waves are $\theta_0 = (k_0 z - \omega_0 t)$ and $\theta_1 = (k_1 z + \omega_1 t)$. We choose

$|\Delta\omega| = |\omega_0 - \omega_1| \ll \omega_0$ and assume a tenuous plasma $\omega_p \ll \omega_{0,1}$, so that $k_0 \approx k_1 \approx \omega_0/c$. By restricting the problem to one dimension we will not address the issue of mutual focusing of counterpropagating lasers [16]. Mutual focusing of copropagating beams, important for high intensity lasers, was described elsewhere [17–19].

Plasma electrons experience a ponderomotive $\vec{v}_1 \times \vec{B}_0 + \vec{v}_0 \times \vec{B}_1$ force, where $\vec{v}_{0,1} = c\vec{a}_{0,1}$ for the lasers of nonrelativistic intensities $a_{0,1} \ll 1$. On a time scale much longer than $1/\omega_0$ electrons are pushed by the intensity gradient of the "optical lattice" and form a density perturbation $\delta n/n_0 = \hat{n} \exp i(\theta_0 + \theta_1) + \text{c.c.}$ Serving as an index grating with the wave number $k = k_0 + k_1$, and oscillating at the difference frequency $\Delta\omega$, this density perturbation scatters the left-going pulse a_1 into the right-going pulse a_0, and vice versa. This is precisely the mechanism of the RBS instability which, in the absence of ion motion, saturates through the particle trapping and wave breaking [20]. It turns out, however, that the superradiant amplification (SRA) of ultrashort pulses is not arrested at the steady state saturation level by particle trapping. Moreover, unlike in the long-pulse experiments, where reflectivity is typically small [21], we find that a superradiant pulse can extract a significant fraction of the pump energy.

To describe the motion of an arbitrary plasma electron (labeled by index j) in the combined field of the two lasers, we introduce a ponderomotive phase $\psi_j = (k_0 + k_1)z_j - \Delta\omega t_j$, where z_j and t_j are the electron position and time, respectively. The equation of motion for the ponderomotive phase, which characterizes the electron position in the optical lattice, is given by

$$\ddot{\psi}_j + \omega_B^2 \sin\psi_j = -\omega_p^2 \sum_{l=1}^{\infty} n_l e^{il\psi_j} - \frac{2\omega_0 e E_z}{mc} + \text{c.c.}, \tag{1}$$

where $\omega_B^2 = 4\omega_0^2 a_0 a_1$ is the bounce frequency of an electron in the optical lattice, and $\hat{n}_l = \langle e^{-il\psi_j}/l\rangle_{\lambda_0/2}$ is the lth harmonic of the small-scale electron plasma wave averaged over one lattice period. The second term in the left-hand side (LHS) of Eq. (1) is the ponderomotive force, the sum in the right-hand side (RHS) is over the harmonics of the spatially varying space-charge force, and E_z is the global electric field. The average of E_z over the lattice period does not vanish, so that it can be thought of as the zeroth harmonic of the short-scale plasma wave. The nonlinear origin of the E_z can be understood as follows. When photons are exchanged between the counterpropagating beams, electrons receive the recoil momentum, producing an average (on $\lambda_0/2$ scale) current. This current in 1D must be balanced by the displacement current [22], producing the electric field E_z which satisfies $\partial E_z/\partial t = -4\pi e\langle J_z\rangle$.

The first harmonic of the plasma density perturbation $\delta n/n_0 = \hat{n}_1 e^{i(\theta_0+\theta_1)} + \text{c.c.}$ causes the PB a_1 to undergo Bragg backscattering into AB a_0, and vice versa. The nonlinear interaction between the beams is calculated by substituting the modified plasma density into the respective eikonal wave equations. For the amplified pulse we obtain

$$2ik_0\left(\frac{\partial}{\partial z} + \frac{1}{c}\frac{\partial}{\partial t}\right)a_0 = -k_p^2 \chi_0 a_0, \tag{2}$$

where $\chi_0 = -\langle e^{-i\psi_j}\rangle a_1^*/a_0$ is the mutual refraction index witnessed by the beam 0. A similar equation is derived for the pump a_1.

The mutual refraction index χ_0 is a complex number. Its real part describes mutual focusing and defocusing of the beams: if $\partial \text{Re}(\chi_0)/\partial r > 0$, beam 0 is focused by beam 1 [16]. The imaginary part of χ_0 describes the energy (or photon) exchange between the two beams: $k_0 \partial|a_0|^2/\partial z = -k_p^2|a_0|^2 \text{Im}(\chi_0)$. If the frequency detuning between the beams $\Delta\omega \approx -\omega_p$, in the linear regime beam 0 is exponentially amplified by the pump at the growth rate $\gamma_{\text{RBS}} = a_1\sqrt{\omega_0/4\omega_p}$ [10,23]. Therefore, linear pulse amplification is readily achieved by precisely detuning the seed signal and the pump. The challenge is, as it turns out, to sustain the short duration of the signal $\tau_L \approx \omega_p^{-1}$ as its intensity grows by several orders of magnitude.

For subrelativistic pump intensities $a_1 \ll 1$, the bandwidth of the RBS (roughly equal to γ_{RBS}) is very narrow because $\gamma_{\text{RBS}} \ll \omega_p$. This implies that the linear amplification suffers from an important limitation: pulses shorter than a plasma period inevitably increase in duration to $\approx \gamma_{\text{RBS}}^{-1}$. Moreover, the amplified pulse lengthens even further as it propagates through the plasma and interacts with the pump because of the effect known as gain narrowing: frequencies in the immediate vicinity of $\omega_1 - \omega_p$ are amplified more than are detuned frequencies [24]. One approach to increasing the amplification bandwidth in the linear regime is to increase the pump intensity. In the so-called strongly coupled regime [23] $\gamma_{\text{RBS}}/\omega_p = \sqrt{3}(\omega_0/16\omega_p)^{1/3} a_1^{2/3} \gg 1$, provided $a_1^2 > 2\omega_p/\omega_0$. The required pump intensity may be too high to make this regime practically interesting. Gain narrowing will still lead to the temporal broadening of the pulse. It turns out, however, that in the nonlinear regime the opposite is true: a short pulse is further shortened in the course of its amplification. Below we demonstrate how *nonlinear* superradiant amplification can be used for efficient amplification of ultrashort pulses.

From Eq. (1), plasma response becomes nonlinear when $\omega_B^2 > \omega_p^2$. In this regime electron bounce motion in the ponderomotive bucket, created by the interference of the two lasers, dominates over the collective interaction with other electrons via the space-charge field of the plasma wave. This is because all the terms in the RHS of Eq. (1) become much smaller than the ponderomotive term in the LHS. Hence, the particle motion is described by the nonlinear oscillator equation $\ddot{\psi}_j + \omega_B^2 \sin\psi_j = 0$. In the reference frame of the beat wave the initially stationary electrons enter the ponderomotive bucket with the initial "speed" $\dot{\psi} = -\Delta\omega$. If this speed is smaller than the bucket height $\dot{\psi}_{\max} = 2\omega_B$, electrons become trapped and execute a synchrotron oscillation in the bucket.

In this strongly nonlinear regime, Eqs. (1) and (2) reduce to the ones analyzed by Bonifacio *et al.* [25] in the context of free-electron lasers (FELs). Those are known to exhibit superradiance: a short intense radiation spike of duration $\tau_L \approx \pi/\omega_B$ is continually amplified. Plasma electrons execute half a bounce in the ponderomotive bucket inside the spike and are left behind the pulse with average negative momentum roughly equal to the bucket height: $\Delta P = -m\omega_B/k_0$. The electron dynamics is crucial for amplification because, when a pump photon backscatters into the pulse, the momentum difference is transferred to the plasma. The ratio of the energy gained by the amplified pulse ΔU to the momentum transferred to the plasma ΔP is given by $\Delta U/\Delta P = -\omega_0/2k_0$. Therefore, efficient and highly localized momentum transfer, which is achieved inside the superradiant spike, is necessary for the likewise amplification of a_0.

A simple qualitative estimate of the pulse evolution can be obtained by assuming that the pulse duration remains $\tau_L \approx \pi/\omega_B$ even as its amplitude grows. Since $\omega_B \propto \sqrt{a_0}$, this implies that the signal slowly narrows as it is amplified. Estimating the total pulse energy (per unit area) as $U = \frac{m^2c^2\omega_0^2}{4\pi e^2}(a_0^2 c\tau_L/2)$ and the electron momentum as $\Delta P = -\Delta z m n_0 \omega_B/k_0$, where Δz is the traveled distance, the energy-momentum balance results in

$$|a_0| \approx \frac{4\omega_p}{3\pi\omega_0} k_p z |a_1| \,. \qquad (3)$$

This result suggests calling the amplified pulse superradiant: intensity of the spike is proportional to $(n_0 z)^2$, i.e., to the total number of plasma electrons encountered by the pulse. This agrees with the classic definition of the superradiance by Dicke [26], according to which the constructive interference of the radiation emitted by a collection of $\mathcal{N}$ independent emitters results in an intensity enhancement which scales as $\mathcal{N}^2$.

Equation (3) is, of course, approximate. The true pulse evolution is unlikely to be exactly superradiant. This is because superradiance requires that the emitters (plasma electrons in our case) not interact with each other in any way other than through the stimulated radiation. In the plasma, in addition to interacting through the ponderomotive force of the laser beat, electrons interact through nonlinear plasma waves, as described by the RHS of Eq. (1). To include these effects we solved Eq. (1) for the plasma and Eq. (2) for the amplified signal a_0 (with a similar equation for the pump a_1) using a 1D time-averaged particle code. To model very long propagation distances, the simulation is done in the reference frame moving with the amplified pulse. This way only a small portion of the entire pump and the plasma, currently located inside the simulation box, is modeled at any instance of time. The simulation box is split into bins, each bin modeling a single ponderomotive bucket. Typical size of the simulation box is $40/k_p$, number of bins is 200, and each bin contains 64 particles which are initially uniformly distributed in the ponderomotive phase ψ_j. Similar codes have been used in the context of free-electron lasers [27]. For a special class of an FEL with an electromagnetic wiggler, the transfer of photons between the wiggler and the amplified signal was studied by Similon and Wurtele [28].

We simulated a collision between a short $\tau_L = \omega_p^{-1}$ signal of initial amplitude $a_0 = 0.025$ and a long pump with $a_1 = 0.025$ in the $n_0 = 2.5 \times 10^{18}$ cm^{-3} plasma. The initial signal amplitude a_0 was chosen to ensure that $\omega_B \approx \omega_p$ at the entrance into the plasma, so that the nonlinear regime is accessed from the start. Taking a smaller initial a_0 resulted in the initial spreading of the pulse, consistently with the linear predictions. The frequency detuning was chosen as $\Delta\omega = -1.7\omega_p$, so that from the very beginning electrons were entering the pulse near the maximum of the ponderomotive bucket. As shown in Fig. 1, the signal is amplified by a factor of 100 in intensity to a mildly relativistic $I = 2 \times 10^{17}$ W/cm^2. Its final FWHM is about one-third of the original. SRA is not very sensitive to the frequency detuning $\Delta\omega$ between the pulse and the pump: in a set of numerical simulations we varied $-2.0\omega_p < \Delta\omega < -1.5\omega_p$ and did not see any qualitative differences in the SRA. Equation (3) predicts that $|a_0/a_1| = 10$ after $k_p z = 500$. From Fig. 1 the ten-fold increase in amplitude occurs at $k_p z = 700$, which is in a qualitative agreement with the fairly crude analytic estimate. Remarkably, energy flows from the pump into the (100 times more intense) amplified pulse. The superradiant signal does not saturate through particle trapping because of its ultrashort duration and because, propagating with the speed close to the speed of light, it constantly encounters unperturbed plasma.

As indicated by Fig. 1, pump depletion is small in tenuous plasmas. This translates into the low efficiency of

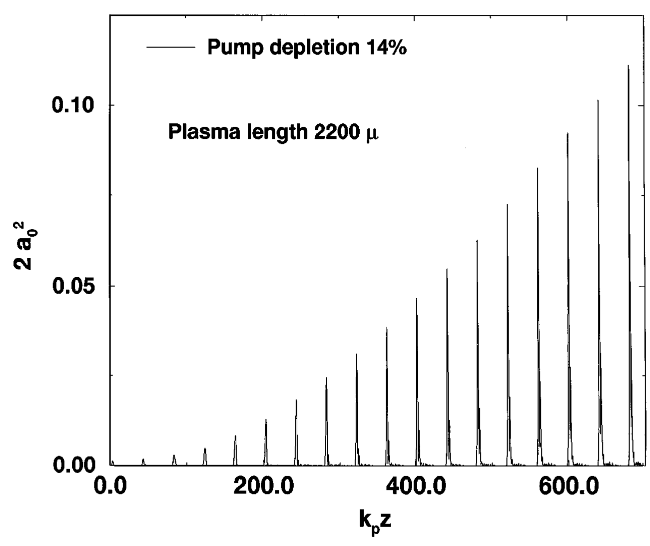

FIG. 1. Superradiant amplification by a factor of 100, results of the time-averaged code. Pump $a_1 = 0.025$; initial signal $a_0 = 0.025$ and $\tau_L = 1/\omega_p$; plasma $n_0 = 2.5 \times 10^{18}$ cm^{-3}.

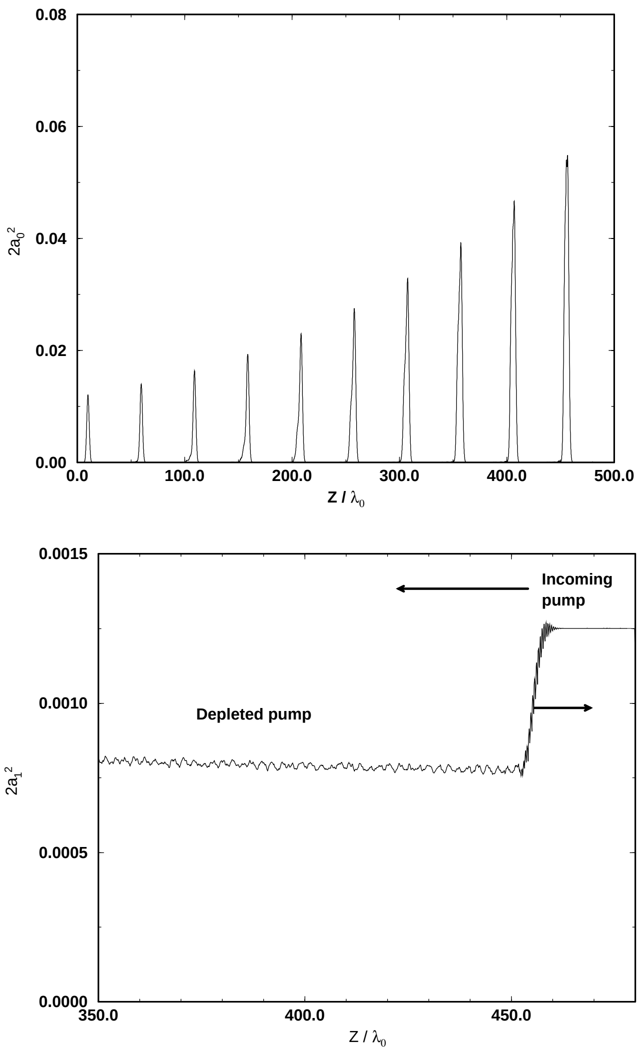

FIG. 2. Superradiant amplification in the large pump depletion regime, results of the PIC simulation with 1D VLPL code. Pump $a_1 = 0.025$; initial signal $a_1 = 0.07$ and $\tau_L = 1/\omega_p$; plasma $n_0 = 10^{19}$ cm^{-3}. Top: Amplification of the short pulse as a function of z. Bottom: Pump depletion.

pulse compression. To demonstrate that much higher extraction efficiency is possible, we simulated a higher density case $\omega_0/\omega_p = 10$ using the one-dimensional version of the VLPL PIC simulation code. As before, the pump amplitude is $a_1 = 0.025$, and the initial signal intensity is $a_0 = 0.07$. The results are shown in Fig. 2. By the time the signal intensity grows by a factor of 5 pump depletion reaches 40%. The intensity amplification and pulse narrowing are less prominent in this example than they are in Fig. 1. This is because, due to the computational demands of the direct PIC simulation, the propagation distance in Fig. 2 is almost 5 times shorter than in Fig. 1.

In conclusion, we demonstrated analytically and numerically the possibility of superradiant amplification of short laser pulses by a properly tuned counterpropagating laser pump. The efficient and dramatic amplification of a sufficiently intense ultrashort pulse by SRA is an approach to the efficient and simple pulse compression.

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