Zettawatt-exawatt lasers and their applications in ultrastrong-field physics

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Since its birth, the laser has been extraordinarily effective in the study and applications of laser-matter interaction at the atomic and molecular level and in the nonlinear optics of the bound electron. In its early life, the laser was associated with the physics of electron volts and of the chemical bond. Over the past fifteen years, however, we have seen a surge in our ability to produce high intensities, 5 to 6 orders of magnitude higher than was possible before. At these intensities, particles, electrons, and protons acquire kinetic energy in the megaelectron-volt range through interaction with intense laser fields. This opens a new age for the laser, the age of nonlinear relativistic optics coupling even with nuclear physics. We suggest a path to reach an extremely high-intensity level 10^{26-28} W/cm^2 in the coming decade, much beyond the current and near future intensity regime 10^{23} W/cm^2, taking advantage of the megajoule laser facilities. Such a laser at extreme high intensity could accelerate particles to frontiers of high energy, teraelectron volt, and petaelectron volt, and would become a tool of fundamental physics encompassing particle physics, gravitational physics, nonlinear field theory, ultrahigh-pressure physics, astrophysics, and cosmology. We focus our attention on high-energy applications, in particular, and the possibility of merged reinforcement of high-energy physics and ultraintense laser.

I. INTRODUCTION

Over the past fifteen years, we have seen a revolution in laser intensities [1]. This revolution stemmed from the technique of chirped pulse amplification (CPA), combined with recent progress in short-pulse generation and superior-energy-storage materials such as Ti:sapphire, Nd:glass, and Yb:glass. The success of this technique was due to its general concept, which fits small, university-type, tabletop-size systems as well as large existing laser chains built for laser fusion in national laboratories such as CEA-Limeil in France, Lawrence Livermore National Laboratory (LLNL), Los Alamos National Laboratory, and the Naval Research Laboratory in the U.S., Rutherford in the U.K., Max Born Institute in Germany, and the Institute of Laser Engineering in Osaka, Japan. A record peak power of petawatt (PW $= 10^{15}$ W) has been produced at LLNL. CPA lasers have given access to a regime of intensities that was not accessible before, opening up a fundamentally new physical domain [2]. The large leap in intensities that we have experienced recently is illustrated in Fig. 1. This figure displays the focused intensities delivered over the years by tabletop systems. After a rapid increase in the 1960s with the invention of lasers, followed by the demonstration of Q switching and mode locking, the power of lasers stagnated due to the inability to amplify ultrashort pulses without causing unwanted nonlinear effects in the optical components. This difficulty was removed with the introduction of the technique of chirped pulse amplification, which took the power of tabletop lasers from

the gigawatt to the terawatt—a jump of 3 to 4 orders of magnitude. This technique was first used with conventional laser amplifiers and more recently extended to optical parametric chirped pulse amplifiers (OPCPA) [3]. A
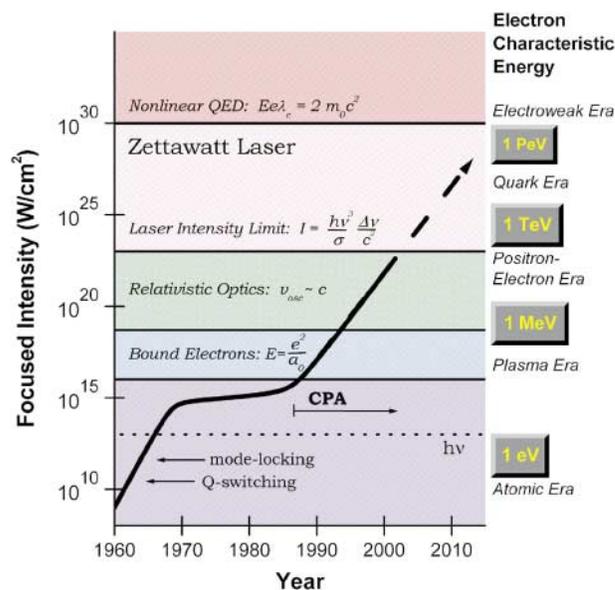

FIG. 1. (Color) Laser-focused intensity vs years for tabletop systems. This system shows the dramatic increase in intensities experienced over the past few years. The figure shows also that the maximum intensity for tabletop systems will be reached in a few years. For the systems we propose, the laser intensity will far exceed these values.

number of laboratories are presently equipped with CPA ultrashort-pulsed terawatt lasers such as Laboratoire Optique Applique in France, University of Lund in Sweden, Max-Planck Institute in Garching, Jena University, and the Japan Atomic Energy Research Institute in Kansai. The CPA-enabled short-pulse generation has advanced to the single-cycle regime [4]. The peak power in the 1990s reached 100 TW, demonstrated at JAERI Kansai. More recently, deformable mirrors have been incorporated into CPA, making it possible with low $f\#$ parabola to focus the laser power on a 1 μm spot size [5]. Present systems deliver focused intensities in the 10^{20} W/cm^2 range. In the near future, CPA systems will be able to produce intensities of the order of 10^{22} W/cm^2. As indicated in Fig. 1, we will see a leveling off of laser intensity for tabletop-size systems at 10^{23} W/cm^2. This limit [1] is imposed, as we will explain later, by the saturation fluence energy per unit area of the amplifying medium and the damage threshold of the optical elements. In a well-conceived CPA system the saturation fluence is of the order but less than the damage threshold. Once this limit is reached, we will have accomplished a leap in intensities of 8 orders of magnitude, and the only way to increase the focused intensity further will be by increasing the beam size, leaving the following questions to be answered: How could we go higher in intensity? And how much higher?

Because the highest intensities will rely on the largest pump available, we explore the possibility whether it could be technically feasible to build a large scale CPA (OPCPA) pumped by a megajoule system of the type of the NIF (National Ignition Facility) in the U.S. and the LMJ (Laser Megajoule) in France. Power in the zettawatt range (10^{21}) could be produced, yielding a focused intensity of 10^{28} W/cm^2. These intensities well beyond the current intensity accessible will open up a new physical regime.

II. THE THEORETICAL INTENSITY LIMIT

In a CPA system the maximum energy per pulse obtainable is limited (1) by the damage threshold F_d of the stretched pulse—of the order of a few nanoseconds, and (2) by the saturation fluence $F_{sat} = h\nu/\sigma$ of the amplifier. Here h is the Planck constant, ν is the laser frequency, and σ is the cross section of the lasing transition at the lasing frequency ν; F_{sat} corresponds to the energy per unit area, necessary to decrease the population inversion by a factor of 2, in a time shorter than the excited state lifetime, T_f. For stretched pulses of the order of a nanosecond, F_d is of the order of 20 J/cm^2 and $F_{sat} = 1$ J/cm^2 for Ti:sapphire. For Yb:glass, $F_{sat} = 40$ J/cm^2—significantly greater than the damage threshold of most optical components. In that regard, high-damage-threshold materials need to be developed in order to take full advantage of the excellent energy storage capability of this material. The minimum pulse duration τ_p is imposed by the Heisenberg uncertainty relation $\Delta\nu_a\tau_p = 1/\pi$, where $\Delta\nu_a$ is the transition bandwidth for

a homogeneously broadened bandwidth. The maximum power per unit area is, therefore, given by

$$P_{th} = \pi \frac{h\nu}{\sigma} \Delta\nu_a. \tag{1}$$

While the maximum focusable intensity I_{th} will be obtained when this power is focused on a spot size limited by the laser wavelength λ, leading to the expression

$$I_{th} = \pi \frac{h\nu^3}{\sigma} \frac{\Delta\nu_a}{c^2}. \tag{2}$$

This intensity limit is represented in Fig. 1.

A. The zettawatt system

Although a zettawatt system could be built using Yb:glass, with the advantages of being relatively compact due to the high F_{sat} of this material and being diode pumpable, much development work needs to be accomplished to reach this intensity level with this material. The proposed systems described below have been stimulated by the construction, both in France and in the U.S., of lasers delivering a few megajoules of energy as well as the availability of large telescope technology (10 m diameter) and deformable mirrors.

Let us recall that the NIF and LMJ systems will deliver 2 MJ at 350 nm, i.e., the third harmonic of 1060 nm in 3 ns. The energy at 530 nm, for a long pulse of 10–20 ns, could be as high as 5 MJ. This energy could be used to pump a CPA-type system to produce of the order of 1 MJ of energy in 10 fs on a spot size of 1 μm [5]. Such a system will have a power of the order of 10^{20} W or 100 exawatts with an intensity of 10^{28} W/cm^2. To produce these phenomenal characteristics, we have two alternatives. The first one would be to use this large pump energy to drive an OPCPA system. This elegant technique [3] has been demonstrated to the joule level at 800 fs a few years ago by Ross *et al.* from Rutherford [6] and is being implemented in the same laboratory to the 10 PW, 30 fs level. It has the advantage that it could use KH$_2$PO$_4$ (KDP) as the nonlinear medium. KDP has been produced in large dimensions at a relatively low cost for the NIF and LMJ programs. The working fluence on the crystal will be 1 J/cm^2, leading to a beam diameter of around 10 m^2 to accommodate the whole pump energy. This technique has the potential to produce a 10-fs pulse or shorter, but needs to be demonstrated at least at the joule level. The second approach, though more conservative, will be to drive a CPA Ti:sapphire, a well-established technology. The amplifier will be composed of a large matrix of Ti:sapphire rods. The size of the matrix will be dictated by the saturation fluence $F_{sat} = 1$ J/cm^2, corresponding to a 10 m diameter beam. The amplifier could be composed of 2500, 20 × 20 cm pieces, each 2 cm in length. Each piece will have to be segmented and isolated in order to avoid the transverse amplified

spontaneous emission. Ti:sapphire of 20×20 cm has already been grown. In both schemes the beams can be focused by a large parabola. The dielectric coating used for the parabola will have a damage threshold of 1 J/cm^2 (for short pulses) [7,8], imposing a parabola size of 10 m in diameter. This parabola will have the same diameter as the Keck telescope. The phase front will be interferometrically controlled by an active matrix of deformable mirrors. The grating compressor could be made out of meter-size gratings assembled interferometrically in a matrix geometry. The size of the grating for a megajoule short-pulse system will also be of the order of 10 m in diameter, dictated by a damage threshold of 1 J/cm^2 (such a damage threshold has just been reported [9]). Each large grating will be composed of 100 of 1 m^2-size gratings.

Note that although large, the number of optical components involved will be small, compared to a MJ system calling for 4500 laser slabs, 800 large KDP frequency converter crystals, and 500 gratings of meter size. Finally, let us mention an elegant [10] alternative method that has been proposed using plasma as a way to transfer the energy from a long pulse into a short pulse. This scheme based on plasma does not suffer from the damage problem. However propagating the laser pulse in a plasma while maintaining good beam quality remains to be demonstrated.

B. The exawatt laser system

If a zettawatt laser, although feasible, could seem too grandiose at this time, an exawatt system, on the other hand, which would produce 10 kJ in 10 fs, i.e., 10^{25} W/cm^2, could be readily constructed. Only 1% or 30 kJ of the NIF/LMJ energy would be necessary. The beam size will be of the order of 1 m in diameter. The amplifying method will be composed of a matrix of 25 Ti:sapphire 20×20 cm^2 crystals and two gratings of meter size. The segmented telescope will have a 1 m aperture. The wave front will be corrected by a large deformable mirror.

In the following we explore a few examples of applications of such intense lasers that may enable new ways to investigate fundamental physics.

III. FUNDAMENTAL PHYSICS: APPLICATIONS OF ZETTAWATT AND EXAWATT LASERS

The major signpost of contemporary high-field science is the entry into the relativistic regime, which is characterized by the quivering momentum of electrons in the laser fields reaching the speed of light times the electron rest mass. This field is reached when the laser intensity is on the order of 10^{18} W/cm^2 for a typical optical frequency of the laser we mentioned above. When the laser field is much less than this, free electrons behave harmonically to the field's optical oscillations. Although these oscillations of electrons are still important enough to couple with various collective motions of free electrons, there is small orbital nonlinearity associated with the quivering motion in this regime. For bound electrons in atoms or molecules the quivering motion in the laser fields is once again, by and large, a perturbation on their orbital motion. This perturbation is enough to cause various interesting phenomena in the nonrelativistic regime. There is such a wealth of phenomena known in this regime that it is hard to list them briefly, but they include the multiphoton process of ionization, that of transitions, Raman and Brillouin scatterings, and various optical nonlinearities [11] arising from the material's response. In the relativistic regime with the intensity above 10^{18} W/cm^2, in addition to the above phenomena, there emerge new classes of effects largely arising from the relativistic nonlinearities of electrons in the high field.

The field intensity in this regime means that the electron momentum in the light is typically $eE_0\omega_0$, which becomes of the order of and exceeds m_0c, where ω_0 is the laser angular frequency, E_0 the laser electric field, and m_0 the electron rest mass. When the electron momentum exceeds m_0c, the electronic orbit ceases to be harmonic and linear. It becomes a figure-eight motion, including higher-harmonic components. The photon pressure is exerted individually on an electron by the electron-photon collision through the Thomson cross section

$$\sigma_T = \frac{8\pi}{3} \frac{e^4}{m_0^2 c^4} \approx 7 \times 10^{-25} \text{ cm}^2. \qquad (3)$$

When the flux of laser at the intensity entering the relativistic regime is shone on an electron (with Lorentz factor γ), this causes a force on it

$$F = \frac{\sigma_T}{\gamma} \frac{E_0^2}{4\pi}. \qquad (4)$$

When the laser intensity is $I = 10^{26}$ W/cm^2, the force acting on an electron is 10^{-2}–10^{-1} erg/cm. Or the acceleration acting on an electron (originally) at rest is

$$a_e = \frac{f}{m_0} \sim 10^{25-26} \text{ cm/s}^2, \qquad (5)$$

where m_0 is the rest mass. (In contrast to this, the Schwinger acceleration is

$$a_S = 2 \times 10^{31} \text{ cm/s}^2, \qquad (6)$$

at which an electron gains energy by m_0c^2 over the Compton length $\lambda_C = \hbar/m_0c$, and the pair creation becomes prevalent.) A large flux of photons bombarding an electron causes such high acceleration through collisions between photons and electrons (Compton collisions). Such acceleration may be called the Eddington acceleration, as Eddington introduced the stellar luminosity at which the gravitational pull is balanced by this photon collisional acceleration (the Eddington luminosity).

The Eddington acceleration may be likened to the acceleration of water molecules near the surface of water in a lake when there is a breeze passing over the water

surface. The collisional viscosity created by collisions between water molecules and wind molecules gives rise to a water flow. However, we also observe that when the breeze gets stronger or becomes a gale, the water surface is no longer smooth and acquires ripples or waves. Such ripples facilitate an increase in the effective viscosity of water for the wind, so that the wind momentum is anomalously effectively transported to water molecules with a much faster rate through the Kelvin-Helmholtz instability at the interface of the two fluids. What happens in the case of strong photon flux ("wind") in a plasma? Just like the strong wind on the water surface, the strong photon flux is capable of creating plasma waves, which in turn causes enhanced viscosity and thus anomalous momentum transport from photons to electrons of the plasma. This latter process is through a collective interaction. The plasma wakefield excitation [12] is typical of this, in which plasma waves generated in this process are accentuated, and new processes of collective interaction emerge. In this ultraintense regime, electrons may be accelerated not only through the electrostatic field that is set up by the ponderomotive force of the laser, but also directly by the ponderomotive force itself to very high energies. In high-intensity regimes, multiphoton processes of electron scattering and acceleration also become important [13]. If we apply this laser at the resonance absorption at high densities near the compressed laser fusion target, on the other hand, much of the laser energy may be converted into relatively low-energy copious electrons, which could constitute a new alternative to the fast ignition fusion. The acceleration of heavier particles (protons and other nuclei) to relativistic energies will become possible, too. Either by direct baryon acceleration by this, or another process (it will take a variety of experimental realizations such as a target irradiation, cluster irradiation, converging imploding shells, etc.), we will access the nuclear regime of matter reminiscent of the early epoch of the big bang (see Fig. 1). The production of extremely high-energy or copious gamma rays will happen. The extreme high photon pressure (which already exceeds a Gbar in the presently available intense lasers) may be finally utilized to directly compress matter in this ultrarelativistic regime, because baryons, too, become relativistic. If so, unprecedented densities of matter may be created. The combination of extreme intensity lasers and high-energy particle beams that can be created by the conventional high-energy physics accelerator will further multiply our ability to expand our frontier horizon. This will be the merging point of high-energy physics and high-field science. In the following, we list several examples of exciting new frontiers of fundamental physics that may be explored by this regime of intensities. In many of the applications we discuss, in this high-intensity regime the interaction length between the laser and matter is expected to be extended beyond the Rayleigh length, as the relativistic mass effect of electrons sets the self-focusing threshold at $10^{10}(\omega/\omega_p)^2$ W [14].

A. Particle acceleration

The pulse or self-modulation of a photon wave packet with sufficient intensity induces a longitudinal electric field (in the x direction, as part of plasma oscillations) as

$$E_x \sim \sqrt{\frac{n_e}{n_{18}}} a_0^2 \text{ (GeV/cm)} \qquad (a_0 \lesssim 1), \qquad (7)$$

where $n_{18} = 10^{18}/\text{cm}^3$ and n_e is the electron density, and $a_0 = eE_0/m_0\omega_0 c$ is the normalized vector potential of the laser, which is sometimes called the quivering velocity normalized to c (or quivering momentum normalized to $m_0 c$). The energy gain over the interaction length l_x is

$$\Delta\varepsilon \sim \sqrt{\frac{n_e}{n_{18}}} a_0^2 \ell_x \text{ (GeV)}. \qquad (8)$$

When a_0 exceeds unity (ultrarelativistic, $I \geq 10^{18}$ W/cm^2), the pressure of the photon wave packet becomes so large that nearly all electrons are evacuated from the laser packet [15]. The photons plow through the plasma with electrons piling up in front of the pulse ("snow plow" of electrons), yielding the snow plow acceleration [16] with momentum gain of electrons

$$p_x = \frac{1}{4}\frac{E_0^2}{n_e c} = \frac{E_c^2}{n_e m_0 c^2} a_0^2, \qquad (9)$$

where $E_c = m_0\omega_0 c/e$. Here the laser pulse is assumed to have fully interacted with the plasma. If the pulse quickly diffracts before the full interaction, say over the Rayleigh length, p_x is simply proportional to E_0.

Now let us imagine, for the moment, that the laser has only half the period (unipolar) [17] or subcyclic [18]. In this case the energy (or momentum) gain is

$$\Delta\varepsilon \sim m_0 c^2 a_0^2. \qquad (10)$$

Since the Lawson-Woodward theorem [19] prohibits any overall acceleration for fully oscillatory (i.e., usual) electromagnetic waves in vacuum in infinite space, the above energy gain is compensated for by the decelerating phase. There are, however, many instances that break the theorem requirements. For example, it may be possible in this extreme relativistic regime that electrons are accelerated to very high energy, immediately reaching the speed of light and becoming in phase with the photon over a sufficiently long distance, so that by the time they become dephased, the electromagnetic wave may decay away for some reason, such as by radiative decay or pump depletion to the acceleration. If this happens, the above energy gain (or a portion of it) may be preserved. In extreme high-intensity regimes such effects will become significant. If the laser spreads over the Rayleigh length, the energy gain is proportional to a_0. The transverse momentum gain (p_y) is always proportional to a_0. In a broad general way we can say that the interaction between the laser and electrons becomes more coherent, as the laser intensity increases because the laser and electrons move more coherently over greater distance. This is

the signature of the acceleration by the ponderomotive potential of photon fields in ultrarelativistic intensity regimes.

Thus it is possible to see electrons at energies of up to ~ 100 TeV at the laser intensity of 10^{26} W/cm^2 and even up to ~ 10 PeV at 10^{28} W/cm^2. The accelerating gradient is 200 TeV/cm and 2 PeV/cm, respectively. Note that such energies (100 TeV and 10 PeV) if collided, correspond to 10^{19} and 10^{23} eV for fixed target experiments. These energies rival or exceed those of the highest energy cosmic rays, which are observed up to 3×10^{20} eV. Of course, it is not easy to attain correspondingly high luminosity for collisions at such high energies. Even though the exploration of particle physics at pb may not be within reach, we may use such particles in pursuit of (other) fundamental physics at the energy frontier. We might recall that Anderson first discovered mesons in cosmic rays, followed by more detailed studies of those particles in cyclotrons and other accelerators. Perhaps the present extreme parameters in the energy frontier may herald some new phenomena. One such example may be the test of Lorentz invariance [20] in extreme high energies. For such a test, unlike the detection of new particles with a pb cross section, the luminosity requirement may be much relaxed.

When we irradiate an extremely relativistic laser pulse ($I \sim 10^{26}$ W/cm^2) on a thin film ($\gtrsim 1$ μm) of a metal at a tight spot of $(1\ \mu\text{m})^2$ followed by a microhole in a metallic slab over more than 1 cm, we hypothesize that the laser pulse picks up metallic electrons from the film and continues to propagate through the microhole, as it remains focused. If this proves to be the case, the amount of electrons that is to be accelerated by this pulse is in the ballpark of

$$N_e \sim n_e A l_b \sim 3 \times 10^{10}, \qquad (11)$$

where we assumed $n_e \sim 10^{24}$ cm^{-3}, area $A = (1\ \mu\text{m})^2$, and the bunch length $l_b \sim 3 \times 10^{-6}$ cm. In such a large pickup, the pump depletion due to the energy transfer to electrons is significant. In fact, it may play a fundamental role in this acceleration to turn the electromagnetic energy into particle kinetic energy without returning to the decelerating phase, providing one way to break the Lawson-Woodward theorem's constraint, as we mentioned earlier. If we take $1/10$ of the above electrons to gain 10 TeV, and if we focus electrons (and positrons) down to 10^{-6} cm (or 10^{-7} cm) at focus (the collision point), the luminosity of the colliding events is of the order of

$$\mathcal{L} = 10^{31} f/\text{cm}^2\,\text{s} \qquad (\text{or } 10^{33} f), \qquad (12)$$

where f is the collision repetition rate. Nakajima has considered a similar but more daring luminosity scenario [21].

B. Fast ignition fusion

One special case of laser energy conversion into electrons is through the resonance absorption at the critical density. The concept of fast ignition in laser-driven inertial fusion [22] calls for a laser beam of ~ 10 psec duration at the intensity exceeding 10^{20} W/cm^2 to be absorbed at the critical density ($\sim 10^{21}$–10^{22}/cm^3), creating a beam of electrons in the several MeV range. The idea is to separate the roles of lasers into two functions: one to compress the fuel with the least amount of entropy increase so that the fusion fuel is compressed to a highest density with the least amount of laser energy, and the other is to heat the fuel to the thermonuclear ignition temperature (~ 10 KeV) when the main compression is achieved. The latter step may be carried out according to Tabak *et al.* by the appropriate range of energetic (several MeV) electrons that are transported from the crust of the target ($<10^{22}$/cm^3) to the surface of the fuel (at $\sim 10^{26}$/cm^3) at the pulse duration of ~ 10 ps. In order for electrons to trigger the fusion ignition, the condition

$$\rho R \lesssim 0.5 \left(\frac{g}{\text{cm}^2}\right) \qquad (13)$$

has to be fulfilled [22]. Here ρ is the density of the compressed fuel and R is the electron range and thus approximately the size of the fuel at compression. Although this idea is potentially capable of reducing the necessary laser energy by almost an order of magnitude or increasing the fusion gain by an order of magnitude at the same laser energy, there remains a considerable uncertainty in the efficiency of energy conversion from a short-pulse (10 ps) laser to the electron beam to the compressed fuel and in the stability and reliability of the electron beam. For example, the laser and the electron beam have to propagate through the over dense plasma, through which filamentation and kinking instabilities are found to arise. In order to cope with this problem, the ignitor laser is further split into the hole boring one and energy deliverer.

We suggest that an alternative method of fast ignition by a much shorter-pulse laser may be possible. In this, although the total amount of energy necessary to be delivered is unchanged (50–100 kJ), we shorten the pulse length to 10 fs so that the local laser intensity reaches of the order of 10^{25} W/cm^2. Since the resonance frequency reduces inversely proportional to $\sqrt{n_e}$, the resonance density becomes on the order of 10^{25}/cm^3, a very close proximity of the fully compressed fuel. This way we may be avoiding the difficult and long energy transport of the electron beam from the density region of $<10^{22}$/cm^2 to 10^{26}/cm^3. With the recent success of improved target design with a conic aperture for fast ignitor laser beam access [23], our ultrafast approach may be further bolstered as a direct energy delivery vehicle. It remains to be seen, however, how much fraction of the laser energy is consumed due to the pump depletion [24] while it interacts with the surrounding plasma. It also needs to be investigated what kind of electron energy spectrum is generated in the ultraintense laser beam. The mission of the electron energy conversion in fast ignition is orthogonal to the previous Sec. III A of super-high-energy electron generation. The production of a small fraction of extremely high-energy electrons is

tolerated, as long as the majority of energy is in several MeV electrons.

C. Baryon acceleration

Many thought it difficult to accelerate protons and heavier particles by light, as massless light propagates at the speed of light, while protons are massive and nonrelativistic—until last year, when the Petawatt Laser experiment [25] and other experiments [26–28] showed that protons have been accelerated much beyond a megaelectron volt. The observed transverse emittance is about 0.5 mm mrad, while the longitudinal one is about MeV psec [29]. These early experiments already rival or even surpass those of the conventional ion sources in some of the crucial parameters. The main mechanism of laser proton acceleration in the above experiments is due to the space charge set up by energetic electrons that are driven forward away from the back surface of the target slab. The energy of protons is thus dependent on that of electrons. Bulanov and others showed [30,31] in simulation that at a laser intensity of $I = 10^{23}$ W/cm^2, protons are accelerated beyond a gigaelectron volt. If this process of proton acceleration scales with the intensity (as the electron energy does), we may be able to see 100 GeV protons and 10 TeV at $I = 10^{26}$ and 10^{28} W/cm^2, respectively. However, it may also be possible that this process is now due directly to the photon pressure beyond the intensity regime of $I = 10^{24}$ W/cm^2. The energy expected through this mechanism is about the same as that through the space charge mechanism.

It is not clear how much energy will be in protons. In the Petawatt Laser experiment, about 10% of laser energy (300 J) was converted into proton energy—30 J (beyond 1 MeV) [25]. If we take this conversion efficiency in the extreme relativistic laser intensity, then more than 1 kJ of proton energy is expected for the case of intensity 10^{26} W/cm^2. If we further take a flat energy spectrum, approximately 10^{11} protons are accelerated beyond 10 GeV in this intensity regime. If we can generate the solitary accelerating structure, such energy may exceed 100 GeV. If we can converge these in a colliding pair of beams at focus, the luminosity of colliding hadrons is $\mathcal{L} = 10^{34} f$, if we focus on 10 nm. The expected number of nuclear events is on the order of 10^9 per shot.

Such an intense, relativistic, compact proton source has a number of fascinating applications. It may be applicable to ion radiography, fast ignition of fusion, etc., among many others. An additional application is pion (or muon) and neutrino beam generation. With sufficiently relativistic proton energies we can speculate that the emittance of created pions is sufficiently small. If so, pions may be promptly accelerated to sufficiently high energies before they die out, before the space charge effect expands the beam emittance. In this the prompt acceleration and its compactness are important, both of which are the forte of the laser acceleration in this application. Protons beyond

a certain energy (several hundred MeV) in matter induce through the nuclear strong interaction the creation of pions, which in turn decay into muons and neutrinos in a matter of 20 ns, if nonrelativistic, propagating at a mere 6 m at most: $\pi \rightarrow \mu + \nu$. With the conventional accelerating gradient of, say, 20 MeV/m over this distance of 6 m, we can increase the pion energy by 120 MeV, which will increase the lifetime of pions but not by an order of magnitude. On the other hand, the laser acceleration with its far-greater gradient would increase the energy and lifetime of pions far more than this (and also reduce the emission cone angle). This would contribute to a further smaller emittance, which in turn contributes to higher energy, lower emittance muons and neutrinos.

D. Gamma ray emission

The emission of gamma rays from the intense laser is expected. Although the well-known bremsstrahlung x rays (and gamma rays) by electrons through the collision with nuclei are expected to remain important, the Larmor radiation is the most intense in the extreme relativistic regime among all radiation mechanisms through the interaction with matter (in this case free electrons) with power

$$P_L = \left(\frac{8}{3} \right) r_e m_0 c \, \omega_0^2 a_0^2. \tag{14}$$

Gamma rays (whose energy peaks at $\hbar\omega_L \sim a_0^3 \hbar\omega_0$) may be forward directed as electrons are accelerated forward. In addition, gamma rays of nuclear origin are also expected. As discussed in Sec. III A, accelerated electrons with an extremely large Lorentz factor created by a laser may be easily converted into extreme high-energy gamma rays (polarization may be preserved if desired).

When an intense laser is directed at a high-energy electron beam, the well-known energy enhancement of the laser photon happens by the factor γ^2 or up to the electron energy itself (where γ is the Lorentz factor of the head-on colliding electron beam) through Compton scattering. Utilizing this, high-energy electron (or positron) colliders may be converted into a γ-γ collider [32]. The idea of a γ-γ collider is an example of the cooperative employment of two separate technologies (the indigenous high-energy physics and the high-field science) to enhance the overall goal.

Gamma rays in the γ-γ collider mode or with an enormous energy as discussed in Sec. III A (with another laser, for example) may be employed to test the Lorentz invariance [20,33]. For example, if Lorentz invariance is violated, the event number of the $e^- e^+$ pair creation is suppressed.

E. Superhot matter

Kishimoto and Tajima [34] have shown that an intense laser pulse may be nearly totally absorbed by only several layers of atomic clusters. When the excursion length

of electrons exceeds the size of clusters, electron orbits become chaotic upon their removal from the original site of the cluster, due to the cluster polarization. The chaos sets in within a few femtoseconds, thus making the absorption of the laser ultrafast. Further, upon removal of electrons, ions of the cluster Coulomb explode, gaining a large fraction of electron energy. If we arrange matter in such a way as to absorb nearly all laser energy over the thickness of a few microns on a $(1 \ \mu m)^2$ spot, the average energy per particle is approximately 100 and $10^4 GeV$, at $I = 10^{26}$ and 10^{28} W/cm^2, respectively. Such superhot matter is expected to generate copious positrons through the Breit-Wheeler process and perhaps other nonlinear quantum electrodynamic (QED) processes [35]. Probably the amount of positrons generated will not be less than 10^9 (see also in Sec. III F).

A large number of nuclear excitations is expected. In fact, in the Petawatt Laser experiment [36] nuclear transmutations through generated gamma rays have been observed. Since the electron energy is much greater than the typical nuclear excitation energy, copious reactions are expected. The nuclear reactivity Y_e is estimated [37] as

$$Y_e = 1.5 \times 10^{12} \frac{n_i n_e \sigma \tau_p \ell}{n_{i0} n_{e0} \sigma_0 \tau_0 \ell_0} \eta \,. \qquad (15)$$

Here n_i and n_e are the ion and electron densities, σ, the reaction cross section, τ_p and ℓ are the interaction time and length, and η is the fraction of reaction electrons. The numbers in the denominator are the normalizing numbers with typical parameters as follows: $n_{i0} = 5 \times 10^{22}/cm^3$, $n_{e0} = 10^{24}$, $\sigma_0 = 1 \ \mu b$, $\tau_0 = 10^{-12}$ s, and $\ell_0 = 10^{-3}$. If we take $\ell \sim 1$ cm, $\eta = 0.5$, and $\sigma \sim 1$ mb, we obtain $10^{14} > Y_e > 10^{10}$. For example, the element of Hf is excited at energy around 1 MeV.

Even if a heavy metal is irradiated, the energy per nucleon exceeds 1 GeV. Thus by irradiating the target from both ends, it may be possible to cause colliding heavy nuclei. The possible nuclear events we expect are

$$N_n = n_i^2 c \sigma V \tau_p \sim 10^7 \,, \qquad (16)$$

where V is the interaction volume taken as 10^{-8} cm^3. We expect about 10^7 events per laser shot of nuclear events, which may include such a process as quark-gluon plasma formation. The presumed transition temperature from a nucleon state to a deconfined quark-gluon plasma state is ~ 200 MeV. The relativistic heavy ion collider (RHIC) delivers 10^9 ion beams (Au) at a luminosity of $2 \times 10^{26}/cm^2$ s, yielding a few events of quark-gluon plasma per second.

F. Nonlinear QED and horizon physics

At the intensity 10^{28} W/cm^2, the electric field is only an order of magnitude less than the Schwinger field as discussed below. At this field, fluctuations in vacuum are polarized by laser to yield copious pairs of real electrons and positrons. In collider physics a similar phenomenon

happens when the so-called Y parameter reaches unity. In reality, even below the Schwinger field, the exponential tail of these fluctuations begins to cause copious pair productions.

Though pair production has been demonstrated at SLAC (E144-experiment) [38] by the interaction of a γ ray with an intense laser at intensities in the range of 10^{18} W/cm^2, the direct production of pairs by a high-intensity laser from vacuum remains elusive. The rule of thumb for threshold of pair production derives from the simple argument that it is the field necessary for a virtual electron to gain an energy $2m_0 c^2$ during its lifetime δt, imposed by the Heisenberg uncertainty principle $\delta t = \hbar/m_0 c^2$, the energy gain length, and $c \delta t$ is the Compton length λ_c. Hence, the breakdown field E_S, the Schwinger field, is $E_S = m_0 c^2/e\lambda_c$ where $\lambda_c = 0.386$ pm $E_S = 2 \times 10^{16}$ V/cm. (The laser field E_S is related to the laser intensity I_ℓ by $E_S^2 = Z_0 I_S$ where Z_0 is the vacuum impedance. For $Z_0 = 377 \ \Omega$, we find a value of $I_S = 10^{30}$ W/cm^2).

This approach gives an estimate for the threshold for pair creation, but does not provide the number of pairs that could be created for a given intensity. The probability of spontaneous production of pair creation per unit time per unit volume by Schwinger [39] is

$$w = \frac{1}{\pi^2} \frac{\alpha}{\delta t} \frac{1}{\lambda_c^3} \left(\frac{E}{E_S}\right)^2 \sum_{n=1}^{\infty} \frac{1}{n^2} \exp\left(-n\pi \frac{E_S}{E}\right), \quad (17)$$

where α is the fine structure constant. The number of pairs N for a given laser field is

$$N = V\tau_p w \,, \qquad (18)$$

where V is the focal volume. For $V = 10^{-12}$ cm^3 and $\tau_p = 10$ fs, we find the generation of 10^{24} pairs at the Schwinger intensity I_S. It is interesting to note that the intensity to create a single pair is still the gargantuan intensity of 10^{27} W/cm^2. However, as indicated earlier (in Sec. III E), the presence of matter such as nuclei or electrons reduces this field by a considerable amount.

The interaction of an intense laser with high-energy electrons will enhance some of the parameters even further. For example, a counterstreaming electron beam and laser can produce copious polarized (high-quality) positrons. This can serve as a polarized positron source for one thing. In an extreme field regime, on the other hand, the laser field is enhanced by the Lorentz factor of the electron beam, so that the effective field from the electron frame may far exceed the Schwinger value. If this field is exceeded by much more than these orders of magnitude, direct production of other particles such as muons out of "vacuum" may be observed.

In addition to the test of nonlinear fields, we are able to explore what may be called "horizon physics." According to Einstein's equivalence principle, a particle that is accelerated feels gravity in the opposite direction of the acceleration. The acceleration due to the electric field of the

laser at this intensity is huge: $a_e \sim 10^{30}$ and 10^{31} cm/s^2, at $I = 10^{26}$ and 10^{28} W/cm^2, respectively. An observer at rest (or in an inertial frame of reference) sees the horizon at infinity if there is no gravitation. On the other hand, an observer near a black hole sees the horizon at a finite distance where the gravitation diverges. Equivalently, an observer who is being accelerated (feeling immense equivalent gravity) now also sees the horizon at a finite distance. Any particle ("observer"— a wave function) that has a finite extent has one side of its wave function leaking out of the horizon. The Unruh radiation is emitted when this happens [40]. Unruh radiation is a sister to the Hawking radiation [41]. The Unruh temperature

$$kT_U = \frac{\hbar a_e}{2\pi c} \qquad (19)$$

which is about 10^4 and 10^5 eV, for $I = 10^{26}$ and 10^{28} W/cm^2, respectively. The radiative power of Unruh radiation increases in proportion to a_0^3 (or $I^{3/2}$), in contrast to the Larmor radiation power of a_0^2 [42]

$$P_U = \frac{12}{\pi} \frac{r_e \hbar}{c} a_0^3 \omega_0^2, \qquad (20)$$

where r_e is the electron classical radius. Since in the extreme relativistic regime the radiation is dominated by Larmor radiation, the Unruh signal has to compete with the Larmor with the ratio of powers

$$\frac{P_U}{P_L} = \frac{\hbar \omega_0}{m_0 c^2} a_0. \qquad (21)$$

At this regime, the Unruh is down only by a few orders of magnitude, and it has been suggested to circumvent the noise to observe the Unruh signal by exploiting polarization, etc. [42]. This P_U yields 10^4 and 10^7 eV/s, for $I = 10^{26}$ and 10^{28} W/cm^2, respectively.

The shrinkage of the distance to the horizon by the violent acceleration allows us to probe other aspects of gravitational physics. For example, Arkani-Hamed *et al.* [43] suggested that extra dimensions of the quantum gravity may have manifestations in a relatively low extra dimension (n) in this four-dimensional world. The distance over which this may be manifested is

$$r_n \sim 10^{32/n-17} \text{ cm}. \qquad (22)$$

The distance to the horizon created by the intense laser acceleration of an electron

$$d = \frac{c^2}{a_e} = \frac{\lambda}{2\pi} \frac{1}{a_0}, \qquad (23)$$

could exceed the above distance Eq. (22) if n is less than or equal to 4.

Zetta- and exawatt lasers will allow us a glimpse into some of the most energetic and enigmatic phenomena of astrophysics such as gamma ray bursts, and their associated effects on extremely high-energy cosmic rays,

the final energy frontier in the Universe ("Extreme Universe" in the highest energy and from the cosmological distance). We marvel at such fascinating possibilities of exploring fundamental physics that are enabled by the extreme high-field science developments by zetta- and exawatt lasers, even though we have just barely scratched the surface so far on this subject. It is also noted that this laser could bring many frontiers of contemporary physics, i.e., particle physics, nuclear physics, gravitational physics, nonlinear field theory, ultrahigh-pressure physics, relativistic plasma and atomic physics, astrophysics, and cosmology together. Because of such significant potential scientific impacts, though it amounts to nontrivial efforts and developments, it seems worthy of further and more serious consideration of this extreme high-field science.

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