Presented at Joint 28th ICFA Advanced Beam Dynamics And Advanced And
Novel Accelerators Workshop On Quantum Aspects Of Beam Physics (QABP03)
7-11 Jan 2003, Higashi Hiroshima, Japan

SLAC-PUB-10426
April 2004

COSMIC PLASMA WAKEFIELD ACCELERATION

PISIN CHEN
Stanford Linear Accelerator Center
Stanford University, Stanford, CA 94309, USA

TOSHIKI TAJIMA
Advanced Photon Research Center
Japan Atomic Energy Research Institute, Kyoto, 619-0215 Japan

YOSHIYUKI TAKAHASHI
Department of Physics, University of Alabama, Huntsville, AL 35899

Recently we proposed a new cosmic acceleration mechanism[1] which was based on the wakefields excited by the Alfven shocks in a relativistically flowing plasma. In this paper we include some omitted details, and show that there exists a threshold condition for transparency below which the accelerating particle is collision-free and suffers little energy loss in the plasma medium. The stochastic encounters of the random accelerating-decelerating phases results in a power-law energy spectrum: $f(\epsilon) \propto 1/\epsilon^2$. As an example, we discuss the possible production of super-GZK ultra high energy cosmic rays (UHECR) in the atmosphere of gamma ray bursts. The estimated event rate in our model agrees with that from UHECR observations.

1. Introduction

Ultra high energy cosmic ray (UHECR) events exceeding the Greisen-Zatsepin-Kuzmin (GZK) cutoff[2] (5×10^{19}eV for protons originated from a distance larger than ~ 50 Mps) have been found in recent years[3,4,5,6]. Observations also indicate a change of the power-law index in the UHECR spectrum (events/energy/area/time $\propto \epsilon^{-\alpha}$) from $\alpha \sim 3$ to a smaller value, at energy around $10^{18} - 10^{19}$eV. These present an acute theoretical challenge regarding their composition as well as their origin[7].

So far the theories that attempt to explain the UHECR can be largely categorized into the "top-down" and the "bottom-up" scenarios. In addition to relying on exotic particle physics beyond the standard model, the main challenges of top-down scenarios are their difficulty in compliance with

Work supported in part by the Department of Energy contract DE-AC03-76SF00515.

the observed event rates and the energy spectrum[7], and the fine-tuning of particle lifetimes. The main challenges of the bottom-up scenarios, on the other hand, are the GZK cutoff, as well as the lack of an efficient acceleration mechanism[7]. To circumvent the GZK limit, several authors propose the "Z-burst" scenario[8] where neutrinos, instead of protons, are the actual messenger across the cosmos. For such a scenario to work, it requires that the original particle, say protons, be several orders of magnitude more energetic than the one eventually reaches the Earth.

Even if the GZK-limit can be circumvented through the Z-burst, the challenge for a viable acceleration mechanism remains, or becomes even more acute. This is mainly because the existing paradigm for cosmic acceleration, namely the Fermi mechanism[9], as well as its variants, such as the diffusive shock acceleration[10], are not effective in reaching ultra high energies[11]. These acceleration mechanisms rely on the random collisions of the high energy particle against magnetic field domains or the shock media, which necessarily induce increasingly more severe energy losses at higher particle energies.

From the experience of terrestrial particle accelerators, we learn that it takes several qualifications for an accelerator to operate effectively. First, the particle should gain energy through the interaction with the longitudinal electric field of a subluminous ($v \leq c$) electromagnetic (EM) wave. In such a setting the accelerated particle can gain energy from the field over a macroscopic distance, much like how a surfer gains momentum from an ocean wave. It is important to note that such a longitudinal field is Lorentz invariant, meaning that the acceleration gradient is independent of the instantaneous energy of the accelerating particle. Second, such a particle-field interaction should be a non-collisional process. This would help to avoid severe energy loss through inelastic scatterings. Third, to avoid excessive synchrotron radiation loss, which scales as particle energy squared, the accelerating particle should avoid any drastic bending beyond certain energy regime. We believe that these qualifications for terrestrial accelerators are also applicable to celestial ones.

Although they are still in the experimental stage, the "plasma wakefield accelerator" concepts[12,13], promise to provide all the conditions stated above. Plasmas are capable of supporting large amplitude electro-static waves with phase velocities near the speed of light. Such collective waves, or "wakefields", can be excited by highly concentrated, relativistic EM energies such as lasers[12] and particle beams[13]. A trailing particle can then gain energy by riding on this wakefield. Although hard scatterings between the accelerating particle and the plasma medium is inevitable, under appropriate conditions, as we will demonstrate below, the particle can be collision-free.

In our recent paper[1] we argued that magneto-shocks (Alfven shocks) in a relativistic plasma flow can also excite large amplitude plasma wakefields, which in turn can be highly efficient in accelerating ultra high energy particles. But with the limited space, many details and intermediate steps were omitted in that paper. Here we provide a more explicit discussion of our notions.

2. Alfven Waves and Plasma Wakefields

It is well-known that an ordinary Alfven wave propagating in a stationary magnetized plasma has a velocity $v_{_A} = eB_0/(4\pi m_i n_p)^{1/2}$, which is typically much less than the speed of light. Here B_0 is the longitudinal magnetic field and n_p is the density of the magnetized plasma. The relative strength between the transverse E and B fields of the Alfven wave is $E/B = v_{_A}/c$. Although the two components are not equal, being mutually perpendicular to the direction of propagation they jointly generate a non-vanishing ponderomotive force that can excite a wakefield in the plasma, which is slow: $v_{ph} = v_A \ll c$. For the purpose of ultra high energy acceleration, such a wakefield would not be too useful, for the accelerating particle can become quickly out of phase with the accelerating field.

Such a slow wave is ordinarily not suitable for accelerating relativistic particles. The situation changes when the plasma as a whole moves with a relativistic bulk velocity $V_p \leq c$. The standard method of obtaining the linear dispersion relation of waves in a magnetized plasma leads to

$$\frac{k_z^2 c^2}{\omega^2} = 1 - \frac{1}{\Gamma_p} \frac{(\omega_{pi}^2 + \omega_{pe}^2)(1 - V_p k/\omega)}{(\omega - V_p k \pm \omega_{Bi}/\Gamma_p)(\omega - V_p k \mp \omega_{Be}/\Gamma_p)} , \tag{1}$$

where k and ω are the wave number and the frequency of the EM wave, respectively, $\omega_{pi,pe} = (4\pi e^2 n_p/m_{i,e})^{1/2}$ are the plasma frequencies for ions and electrons, and $\omega_{Bi,Be} = (eB_0/m_{i,e})^{1/2}$ are the ion and electron cyclotron frequencies. Here Γ_p is the Lorentz factor of the bulk plasma flow. Figure 1 shows the dispersion relations of various transverse EM waves that propagate along the direction of B_0 with and without the plasma bulk flow V_p. In Fig. 1(a) we see that outside the lightcone (superluminous, or $v_{ph} > c$) lie the regular EM waves, whose asymptotic dispersion is $\omega = kc$. Within the lightcone (subluminous), there are two additional branches, the whistler wave (an electron branch mode) and the Alfven wave whose frequency remains quite low and its electric field is much smaller than the magnetic one, i.e., $E/B = v_A/c \ll 1$ in the absence of flow.

In the case where the bulk flow of the plasma approaches the speed of light, however, the Alfven waves acquire a phase velocity close to c and enhances the ratio of E/B to $\sim V_p/c \leq 1$, and it becomes indistinguishable from a bona fide EM wave. Preliminary results from simulations indicate that such relativistic Alfven waves can indeed excite plasma wakefields[14] Further simulation works are currently in progress, as reported in this workshop[15]. In this relativistic flow the excited wakefields are all in one direction, which contributes to the unidirectional acceleration. With our applications to astrophysical problems in mind, the Alfven-wave-plasma interaction relevant to us is in the nonlinear regime.

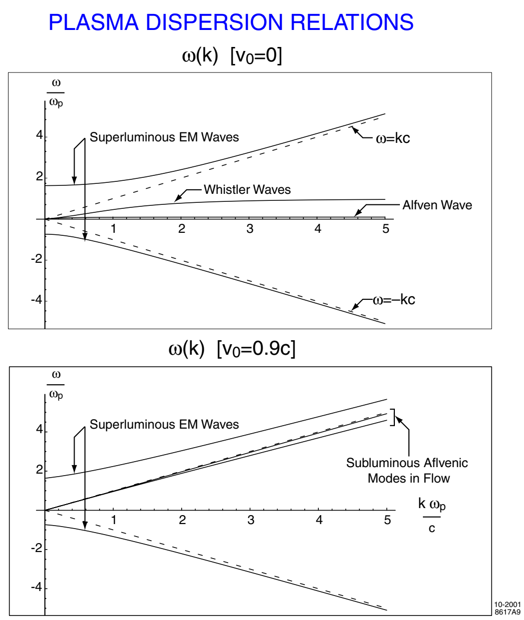

Figure 1. The dispersion relations for stationary and relativistic plasma flows.

The plasma wakefield in the nonlinear regime has been well-studied[16]. The nonlinearity is determined by the driving EM wave's *ponderomotive* potential, which is governed by its normalized vector potential $a_0 = eE/mc\omega$. When this parameter exceeds unity, nonlinearity is strong[12] so that additional important physics incurs. For a stationary plasma, the maximum field amplitude that the plasma wakefield can support is

$$E_{\rm max} \approx E_{\rm wb} a_0 = \frac{m_e c \omega_p}{e} a_0 \ , \tag{2}$$

which is enhanced by a factor a_0 from the cold wavebreaking limit (the naively assumed maximum field), $E_{\rm wb} = m_e c\omega_p/e$, of the linear regime. In a relativistic plasma flow with a Lorentz factor Γ_p, the cold wavebreaking field is reduced by a factor $\Gamma_p^{1/2}$ due to Lorentz contraction. The maximum "acceleration gradient" G experienced by a singly-charge particle riding on this plasma wakefield is then

$$G = eE'_{\rm max} \approx a_0 m_e c^2 \sqrt{\frac{4\pi r_e n_p}{\Gamma_p}} \ . \tag{3}$$

The plasma wavelength, in the mean time, is stretched also by a factor a_0 from that in the linear regime. So in a plasma flow the wavelength is

$$\lambda_{pN} = \frac{2}{\pi} a_0 \lambda'_p \approx a_0 \sqrt{\frac{\pi \Gamma_p}{r_e n_p}} \ , \tag{4}$$

where $r_e = e^2/m_e c^2 = 2.8 \times 10^{-13}$cm is the classical electron radius.

3. Maximum Energy Gain and Spectrum

To determine the maximum possible energy gain, we need to know how far can a test particle be accelerated. At ultra high energies once the test particle encounters a hard scattering or bending, the hard-earned kinetic energy would most likely be lost. The scattering of an ultra high energy proton with the background plasma is dominated by the proton-proton collision. Existing laboratory measurements of the total *pp* cross section scales roughly as $\sigma_{pp} = \sigma_0 \cdot \{1 + 6.30 \times 10^{-3} [\log(s)]^{2.1}\}$, where $\sigma_0 \approx 32$mb and the center-of-mass energy-squared, s, is given in $({\rm GeV})^2$. In our system, even though the UHE protons are in the ZeV regime, the center-of-mass energy of such a proton colliding with a comoving background plasma proton is in the TeV range, so it is safe to ignore the logarithmic dependence and assume a constant total cross section, $\sigma_{pp} \sim \sigma_0 \sim 30$ mb in the ZeV energy regime. Since in astrophysical settings an out-bursting relativistic plasma dilutes as it expands radially, its density scales as $n_p(r) = n_{p0}(R_0/r)^2$, where n_{p0} is the plasma density at a reference radius R_0 . The proton mean-free-path can be determined by integrating the collision probability up to unity,

$$1 = \int_{R_0}^{R_0+L_{\rm mfp}} \frac{\sigma_{pp} n_p(r)}{\Gamma_p} dr = \int_{R_0}^{R_0+L_{\rm mfp}} \frac{\sigma_{pp} n_{p0}}{\Gamma_p} \frac{R_0^2}{r^2} dr \ . \tag{5}$$

We find

$$1 = \frac{\sigma_{pp} n_{p0} R_0}{\Gamma_p} \left[1 - \frac{R_0}{R_0 + L_{\rm mfp}} \right] . \tag{6}$$

Since $L_{\rm mfp}$ is positive definite, $0 < [1 - R_0/(R_0 + L_{\rm mfp})] < 1$. Therefore the solution to $L_{\rm mfp}$ does not exist unless the coefficient, $\sigma_{pp} n_{p0} R_0/\Gamma_p > 1$. That is there exists a threshold condition below which the system is collision-free:

$$\frac{\sigma_{pp} n_{p0} R_0}{\Gamma_p} = 1 . \tag{7}$$

When a system is below this threshold, a test particle can in principle be accelerated unbound. In practice, of course, other secondary physical effects would eventually intervene.

In a terrestrial accelerator, the wakefields are coherently excited by the driving beam, and the accelerating particle would ride on the same wave crest over a macroscopic distance. There the aim is to produce near-monoenergetic final energies (and tight phase-space) for high energy physics and other applications. In astrophysical settings, however, the drivers, such as the Alfven shocks, will not be so organized. A test particle would then face random encounters of accelerating and decelerating phases of the plasma wakefields excited by Alfven shocks.

The stochastic process of the random acceleration-deceleration can be described by the distribution function $f(\epsilon, t)$ governed by the Chapman-Kolmogorov equation[17,18]

$$\begin{aligned} \frac{\partial f}{\partial t} = & \int_{-\infty}^{+\infty} d(\Delta\epsilon) W(\epsilon - \Delta\epsilon, \Delta\epsilon) f(\epsilon - \Delta\epsilon, t) - \int_{-\infty}^{+\infty} d(\Delta\epsilon) W(\epsilon, \Delta\epsilon) f(\epsilon, t) \\ & - \nu(\epsilon) f(\epsilon, t) . \end{aligned} \tag{8}$$

The first term governs the probability per unit time of a particle "sinking" into energy ϵ from an initial energy $\epsilon - \Delta\epsilon$ while the second term that "leaking" out from ϵ. The last term governs the dissipation due to collision or radiation, or both. As we will demonstrate later, the astrophysical environment that we invoke for the production of UHECR is below the collision threshold condition, and so accelerating particles are essentially collision-free.

The radiation loss in our system is also negligible. As discussed earlier, in a relativistic flow the transverse E and B fields associated with the Alfven shock are near equal in magnitude. Analogous to that in an ordinary EM wave, an ultra relativistic particle (with a Lorentz factor γ) co-moving with such a wave will experience a much suppressed bending field, by a factor $1/\gamma^2$. Furthermore, the plasma wakefield acceleration takes place in the region that trails behind the shock (and not in the bulk of the shock) where the accelerating particle in effect sees only the longitudinal *electrostatic* field colinear to the particle motion[16]. We are therefore safe to ignore the radiation loss entirely as well. We can thus ignore the dissipation term in the Chapman-Komogorov equation and focus only on the purely random plasma wakefield acceleration-deceleration.

Assuming that the energy gain per phase encounter is much less than the final energy, i.e., $\Delta\epsilon \ll \epsilon$, we Taylor-expand $W(\epsilon - \Delta\epsilon, \Delta\epsilon) f(\epsilon - \Delta\epsilon)$ around $W(\epsilon, \Delta\epsilon) f(\epsilon)$ in the sink term and reduce Eq.(9) to the Fokker-Planck equation

$$\begin{aligned}\frac{\partial f}{\partial t} = \frac{\partial}{\partial \epsilon}\int_{-\infty}^{+\infty} d(\Delta\epsilon)\Delta\epsilon W(\epsilon, \Delta\epsilon) f(\epsilon, t) \\ + \frac{\partial^2}{\partial \epsilon^2}\int_{-\infty}^{+\infty} d(\Delta\epsilon)\frac{\Delta\epsilon^2}{2} W(\epsilon, \Delta\epsilon) f(\epsilon, t) \ .\end{aligned} \tag{9}$$

We now assume the following properties of the transition rate $W(\epsilon, \Delta\epsilon)$ for a purely stochastic process:

a) W is an even function;

b) W is independent of ϵ;

c) W is independent of $\Delta\epsilon$.

Property a) follows from the fact that in a plasma wave there is an equal probability of gaining and losing energy. In addition, since the wakefield amplitude is Lorentz invariant, the chance of gaining a given amount of energy, $\Delta\epsilon$, is independent of the particle energy ϵ. Finally, under a purely stochastic white noise, the chance of gaining or losing any amount of energy is the same. Based on these arguments we deduce that

$$W(\epsilon, \Delta\epsilon) = \frac{1}{2c\tau^2 G} \ , \tag{10}$$

where τ is the typical time of interaction between the test particle and the random waves and G is the maximum acceleration gradient (cf. Eq.(4)). We note that there is a stark departure of the functional dependence of W in our theory from that in Fermi's mechanism, in which the energy gain $\Delta\epsilon$ per encounter scales linearly and quadratically in ϵ for the first-order and second-order Fermi mechanism, respectively.

To look for a stationary distribution, we put $\partial f/\partial t = 0$. Since W is an even function, the first term on the RHS in Eq.(10) vanishes. To ensure the positivity of particle energies before and after each encounter, the integration limits are reduced from $(-\infty, +\infty)$ to $[-\epsilon, +\epsilon]$, and we have

$$\frac{\partial^2}{\partial \epsilon^2} \int_{-\epsilon}^{+\epsilon} d(\Delta\epsilon) \frac{\Delta\epsilon^2}{2} W(\epsilon, \Delta\epsilon) f(\epsilon) = 0 \ . \tag{11}$$

Inserting W from Eq.(11), we arrive at the energy distribution function that follows power-law scaling,

$$f(\epsilon) = \frac{\epsilon_0}{\epsilon^2} \ , \tag{12}$$

where the normalization factor ϵ_0 is taken to be the mean energy of the background plasma proton, $\epsilon_0 \sim \Gamma_p m_p c^2$. The actually observed UHECR spectrum is expected to be degraded somewhat from the above idealized, theoretical power-law index, $\alpha = 2$, not only due to possible departure of the reality from the idealized model, but also due to additional intermediate cascade processes that transcend the original UHE protons to the observed UHECRs.

We note that a power-law energy spectrum is generic to all purely stochastic, collisionless acceleration processes. This is why both the first and the second order Fermi mechanisms also predict power-law spectrum, if the energy losses, e.g., through inelastic scattering and radiation (which are severe at ultra high energies), are ignored. The difference is that in the Fermi mechanism the stochasticity is due to random collisions of the test particle against magnetic walls or the shock medium, which necessarily induce *reorientation* of the momentum vector of the test particle after every diffusive encounter, and therefore should trigger inevitable radiation loss at high energies. The stochasticity in our mechanism is due instead to the random encounters of the test particle with different accelerating-decelerating *phases*. As we mentioned earlier, the phase vector of the wakefields created by the Alfven shocks in the relativistic flow is nearly unidirectional. The particle's momentum vector, therefore, never changes its direction but only magnitude, and is therefore radiation free in the energy regime that we consider for proton acceleration.

4. Gamma Ray Bursts and Wakefield Acceleration

We now apply our acceleration mechanism to the problem of UHECR. GRBs are by far the most violent release of energy in the universe, second only to the big bang itself. Within seconds (for short bursts) about $\epsilon_{\rm GRB} \sim 10^{52}$erg of energy is released through gamma rays with a spectrum that peaks around several hundred keV. Existing models for GRB, such as the relativistic fireball model[19], typically assume neutron-star-neutron-star (NS-NS) coalescence as the progenitor. Neutron stars are known to be compact ($R_{\rm NS} \sim O(10)$km) and carrying intense surface magnetic fields ($B_{\rm NS} \sim 10^{12}$G). Several generic properties are assumed when such compact objects collide. First, the collision creates sequence of strong magneto-shocks (Alfven shocks). Second, the tremendous release of energy creates a highly relativistic out-bursting fireball, most likely in the form of a plasma.

The fact that the GRB prompt (photon) signals arrive within a brief time-window implies that there must exists a threshold condition in the GRB atmosphere where the plasma becomes optically transparent beyond some radius R_0 from the NS-NS epicenter. Applying Eq.(8) to the case of out-bursting GRB photons, this condition means

$$\frac{\sigma_c n_{p0} R_0}{\Gamma_p} = 1 \ , \tag{13}$$

where $\sigma_c = (\pi r_e^2)(m_e/\omega_{\rm GRB})[\log(2\omega_{\rm GRB}/m_e) + 1/2] \approx 2 \times 10^{-25}{\rm cm}^2$ is the Compton scattering cross section. Since $\sigma_{pp} < \sigma_c$, the UHECRs are also collision-free in the same environment. There is clearly a large parameter space where this condition is satisfied. To narrow down our further discussion, it is not unreasonable to assume that $R_0 \sim O(10^4)$km. A set of self-consistent parameters can then be chosen: $n_{p0} \sim 10^{20}{\rm cm}^{-3}, \Gamma_p \sim 10^4$, and $\epsilon_0 \sim 10^{13}{\rm eV} \equiv \epsilon_{13}$.

To estimate the plasma wakefield acceleration gradient, we first derive the value for the a_0 parameter. We believe that the megneto-shocks constitute a substantial fraction, say $\eta_a \sim 10^{-2}$, of the total energy released from the GRB progenitor. The energy Alfven shocks carry is therefore $\epsilon_A \sim 10^{50}$erg. Due to the pressure gradient along the radial direction, the magnetic fields in Alfven shocks that propagate outward from the epicenter will develop sharp discontinuities and be compactified[20]. The estimated shock thickness is $\sim O(1)$m at $R_0 \sim O(10^4)$km. From this and ϵ_A one can deduce the magnetic field strength in the Alfven shocks at R_0, which gives $B_A \sim 10^{10}$G. This leads to $a_0 = eE_A/mc\omega_A \sim 10^9$. Under these assumptions, the acceleration gradient G (*cf.* Eq.(4)) is as large as

$$G \sim a_0 mc^2 \sqrt{\frac{4\pi r_e}{\sigma_c R_0}} \sim 10^{16} \Big(\frac{a_0}{10^9}\Big) \Big(\frac{10^9{\rm cm}}{R_0}\Big)^{1/2} {\rm eV/cm} \ . \tag{14}$$

Although the UHE protons can in principle be accelerated unbound in our system, the ultimate maximum reachable energy is determined by the conservation of energy and our assumption on the population of UHE protons. Since it is known that the coupling between the ponderomotive potential of the EM wave and the plasma wakefield is efficient, we assume that the Alfven shock energy is entirely loaded to the plasma wakefields after propagating through the plasma. Furthermore, we assume that the energy in the plasma wakefield is entirely reloaded to the UHE protons through the stochastic process. Thus the highest possible UHE proton energy can be determined by energy conservation

$$\eta_a \epsilon_{\rm GRB} \sim \epsilon_{\rm A} \sim \epsilon_{\rm UHE} \sim N_{\rm UHE} \int_{\epsilon_{13}}^{\epsilon_m} \epsilon f(\epsilon) d\epsilon \ . \tag{15}$$

which gives

$$\epsilon_m = \epsilon_{13} \exp(\eta_a \epsilon_{\rm GRB} / N_{\rm UHE} \epsilon_{13}) \ . \tag{16}$$

This provides a relationship between the maximum possible energy, ϵ_m, and the UHE proton population, $N_{\rm UHE}$. We assume that $\eta_b \sim 10^{-2}$ of the GRB energy is consumed to create the bulk plasma flow, i.e., $\eta_b \epsilon_{\rm GRB} \sim N_p \Gamma_p m_p c^2 \sim N_p \epsilon_{13}$, where N_p is the total number of plasma protons. We further assume that $\eta_c \sim 10^{-2}$ of the plasma protons are trapped and accelerated to UHE, i.e., $N_{\rm UHE} \sim \eta_c N_p$. Then we find $\epsilon_m \sim \epsilon_{13} \exp(\eta_a / \eta_b \eta_c)$. We note that this estimate of ϵ_m is exponentially sensitive to the ratio of several efficiencies, and therefore should be handled with caution. If the values are indeed as we have assumed, $\eta_a / \eta_b \eta_c \sim O(10^2)$, then ϵ_m is effectively unbound until additional limiting physics enters. Whereas if the ratio is $\sim O(10)$ instead, the UHE cannot even reach the ZeV regime. The validity of our assumed GRB efficiencies then relies on the consistency check against observations.

5. UHECR Event Rate

In addition to the energy production issue, equally important to a viable UHECR model is the theoretical estimate of the UHECR event rates. The NS-NS coalescence rate is believed to be about 10 events per day in the entire Universe[21,22]. This frequency is consistent with the observed GRB events, which is on the order of $f_{\rm GRB} \sim 10^{3.5}$ per year.

In the Z-burst scenario an initial neutrino energy above 10^{21}eV[8] or 10^{23}eV[23] is required (depending on the assumption of the neutrino mass) to reach the Z-boson threshold. For the sake of discussion, we shall take the necessary neutrino energy as $\epsilon_\nu > 10^{22}$eV. Such ultra high energy neutrinos can in principle be produced through the collisions of UHE protons with the GRB background protons: $pp \to \pi + X \to \mu + \nu + X$. All UHE protons with energy $\epsilon_{>22} \geq 10^{22}$eV should be able to produce such neutrinos. The mean energy (by integrating over the distribution function $f(\epsilon)$) of these protons is $\langle\epsilon_{>22}\rangle \sim O(100)\epsilon_{22}$. Therefore the multiplicity of neutrinos per UHE proton is around $\mu_{(p\to\nu)} \sim O(10) - O(100)$. At the opposite end of the cosmic process, we also expect multiple hadrons produced in a Z-burst. The average number of protons that Z-boson produces is ~ 2.7[24]. Finally, the population of UHE protons above 10^{22}eV is related to the total UHE population by $N_{>22} \sim (\epsilon_{13}/\epsilon_{22})N_{\mathrm{UHE}} \sim \eta_b\eta_c\epsilon_{\mathrm{GRB}}/\epsilon_{22}$.

Putting the above arguments together, we arrive at our theoretical estimate of the expected UHECR event rate on earth,

$$\begin{aligned} N_{\mathrm{UHECR}}(>10^{20}\mathrm{eV}) &= f_{\mathrm{GRB}}\mu_{(p\to\nu)}\mu_{(Z\to p)}N_{>22}\frac{1}{4\pi R_{\mathrm{GRB}}^2} \\ &\sim f_{\mathrm{GRB}}\mu_{(p\to\nu)}\mu_{(Z\to p)}\eta_b\eta_c\frac{\epsilon_{\mathrm{GRB}}}{\epsilon_{22}}\frac{1}{4\pi R_{\mathrm{GRB}}^2}\ . \end{aligned} \tag{17}$$

The typical observed GRB events is at a redshift $z \sim O(1)$, or a distance $R_{\mathrm{GRB}} \sim 10^{23}$km. Our estimate of observable UHECR event rate is therefore

$$N_{\mathrm{UHECR}}(>10^{20}\mathrm{eV}) = O(1)/100\mathrm{km}^2/\mathrm{yr}/\mathrm{sr}\ , \tag{18}$$

which is consistent with observations, or in turn this observed event rate can serve as a constraint on the various assumptions of our specific GRB model.

6. A Laboratory Astrophysics Experiment

History has shown that the symbiosis between direct observation and laboratory investigation was instrumental in the progress of astrophysics. Our cosmic plasma wakefield acceleration mechanism can in principle be tested in the laboratory setting[26]. A schematic diagram for such an experiment is shown in Figure 2.

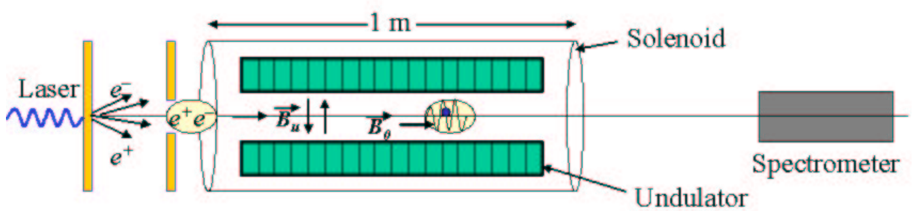

Figure 2. A schematic diagram of a possible laboratory experiment to verify the Alfven-induced plasma wakefield acceleration mechanism.

The main goals for such an experiment are

1. Generation of Alfven waves in a relativistic plasma flow;
2. Inducing high gradient nonlinear plasma wakefields;
3. Acceleration and deceleration of trapped e^+/e^-;
4. Power-law $(n\ -2)$ spectrum due to stochastic acceleration.

Although it is unlikely that the extremely high density, high intensity and high acceleration gradient involved in this acceleration mechanism can be reproduced in the laboratory setting, it is hoped that the key elements necessary for this mechanism can indeed be verified. In this regard, the value of the experiment lies in its validation of the underlying dynamics of the Alfven-induced plasma wakefield acceleration.

7. Accknowledgement

We appreciate helpful discussions with J. Arons, R. Blandford, P. Meszaros. This work was supported by Department of Energy contracts DE–AC03–76SF00515 and DE-AC03-78SF00098.

References

1. P. Chen, T. Tajima, Y. Takahashi, *Phys. Rev. Lett.* **89**, 161101 (2002).
2. K. Greisen, *Phys. Rev. Lett.* **16**, 748 (1966); G. T. Zatsepin and V. A. Kuzmin, *Pis'ma Zh. Eksp. Teor. Fiz.* **4**, 114 (1966) [*JETP Lett.* **4**, 78 (1966).
3. D. J. Bird et al., *Phys. Rev. Lett.* **71**, 3401 (1993); *Astrophys. J.* **424**, 491 (1994); **441**, 144 (1995).
4. M. Takeda et al., *Phys. Rev. Lett.* **81**, 1163 (1998); *Astrophys. J.* **522**, 225 (1999).
5. T. Abu-Zayyad et al., *Int. Cosmic Ray Conf.* **3**, 264 (1999).
6. M. A. Lawrence, R. J. O. Reid, and A. A. Watson, *J. Phys. G Nucl. Part. Phys.* **17**, 773 (1991).
7. A. V. Olinto, *Phys. Rep.* **333-334**, 329 (2000).
8. T. Weiler, *Astropart. Phys.* **11**, 303 (1999); D. Fargion, B. Mele, and A. Salis, *Astrophys. J.* **517**, 725 (1999).
9. E. Fermi, *Phys. Rev.* **75**, 1169 (1949); *Astrophys. J* **119**, 1 (1954).
10. W. I. Axford, E. Leer, and G. Skadron, in *Proc. 15th Int. Cosmic Ray Conf. (Plovdic)* **11**, 132 (1977); G. F. Krymsky, *Dokl. Acad. Nauk. SSR* **234**, 1306 (1977); A. R. Bell, *Mon. Not. R. Astro. Soc.* **182**, 147 (1978); R. D. Blandford and J. F. Ostriker, *Astrophys. J. Lett.* **221**, L29 (1978).
11. A. Achterberg, in *Highly Energetic Physical Processes and Mechanisms for Emission from Astrophysical Plasmas*, IAU Symposium, vol. 195, P. C. H. Martens and S. Tsuruta, eds. (1999).
12. T. Tajima and J. M. Dawson, *Phys. Rev. Lett.* **43**, 267 (1979).
13. P. Chen, J. M. Dawson, R. Huff, and T. Katsouleas, *Phys. Rev. Lett.* **54**, 693 (1985).

14. P. Romenesko, P. Chen, T. Tajima, in preparation (2002).
15. K. Reil, presented in this workshop and in these proceedings.
16. E. Esarey, P. Sprangle, J. Krall, and A. Ting, *IEEE Trans. Plasma Sci.* **24**, 252 (1996).
17. K. Mima, W. Horton, T. Tajima, and A. Hasegawa, in Proc. *Nonlinear Dynamics and Particle Acceleration*, eds. Y.H.Ichikawa and T.Tajima (AIP, New York, 1991) p.27.
18. Y. Takahashi, L. Hillman, T. Tajima, in *High Field Science*, eds. T. Tajima, K. Mima, H. Baldis, (Kluwer Academic/Plenum 2000) p.171.
19. M. J. Rees and P. Meszaros, *Mon. Not. R. Astro. Soc.* **158**, P41 (1992); P. Mezsaros and M. J. Rees, *Astrophys. J.* **405**, 278 (1993).
20. A. Jeffrey and T. Taniuti, *Nonlinear Wave Propagation*, (Academic Press, NY, 1964).
21. T. Piran, ApJ Letters 389, L45 (1992); A. Shemi and T. Piran, *Atrophys. J.* **365**, L55 (1990).
22. LISA collaboration, http://www.lisa.uni-hannover.de/
23. G. Gelmini and G. Varieschi, UCLA/02/TEP/4 (2002), unpublished.
24. Particle Data Group, Phys. Rev. D54, 187 (1996).
25. M. Punch et al., Nature 358, 477 (1992); J. Quinn et al., ApJ Letters 456, L63 (1996); C. M. Urry, *Advances in Space Research* **21**, 89 (1998).
26. P. Chen, Asso. Asian Pacific Phys. Soc. (AAPPS) Bull. **13**, 3 (2003).