Astrophysical ZeV acceleration in the relativistic jet from an accreting supermassive blackhole

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A B S T R A C T

An accreting supermassive blackhole, the central engine of active galactic nucleus (AGN), is capable of exciting extreme amplitude Alfven waves whose wavelength (wave packet) size is characterized by its clumpiness. The ponderomotive force and wakefield are driven by these Alfven waves propagating in the AGN (blazar) jet, and accelerate protons/nuclei to extreme energies beyond Zetta-electron volt (ZeV = 10^21 eV). Such acceleration is prompt, localized, and does not suffer from the multiple scattering/bending enveloped in the Fermi acceleration that causes excessive synchrotron radiation loss beyond 10^{19} eV. The production rate of ZeV cosmic rays is found to be consistent with the observed gamma-ray luminosity function of blazars and their time variabilities.

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1. Introduction

The origin of ultra-high energy cosmic rays (UHECRs) with energies 10^{19} eV remains a puzzle of astrophysics. It is generally believed to be extragalactic ([1], references therein). The production of UHECRs has been discussed mainly in the framework of the Fermi acceleration [2], in which charged particles gain energy through a numerous number of scatterings by the magnetic clouds. One of the necessary conditions of Fermi acceleration is the magnetic confinement: the Hillas criterion sets a constraint on the product of the magnetic field strength B and extension R of the candidate objects (Hillas criterion) [3]: W \leq W_{\text{max}} \sim z(B/1 \mu G)(R/1 \text{kpc}) \text{eV}, where z is the charge of the particle. The possible candidate objects (but only marginally satisfying the Hillas criterion for 10^{19} eV production) are neutron stars, active galactic nuclei (AGN), gamma-ray bursts (GRBs), and accretion shocks in the intergalactic space. However, the acceleration of 10^{20} eV particles even in those candidate objects is not easy for the Fermi mechanism because of (1) the large number of scatterings necessary to reach highest energies, (2) energy losses through the synchrotron emission at the bending associated with scatterings, and (3) difficulty in the escape of particles which are initially magnetically confined in the acceleration domain [1].

In the present paper we point out that there is an alternative way to accelerate charged particles (protons, ions, and electrons) to ultra-high energies in cosmic conditions, in particular in the conditions of AGN, through the electromagnetic (EM) wave-particle interaction. Along this path two conditions are necessary: (a) the accelerating structure (wave) should have a relativistic propagation velocity (phase velocity) very close to the speed of light c; (b) the wave should have a relativistic amplitude (i.e. so large an amplitude that the particle acquires relativistic momentum in one oscillation period of the wave, eE/\omega > mc, where E and \omega are the wave electric field and frequency, e and m are the charge and mass of the jth particle). The condition (b) is needed because the electromagnetic field acceleration can yield acceleration in the direction of the wave propagation only from the nonlinear force of \nu \times B/c, called the ponderomotive force, and this term becomes significant only when the amplitude becomes relativistic [4]. We note that these two conditions may be fulfilled in a number of astrophysical settings (as well as in many modern terrestrial laboratories [5]). When the conditions (a) and (b) are fulfilled, this acceleration mechanism for UHECR generation has advantages over the Fermi mechanism, for the following reasons:

1. The ponderomotive field provides an extremely high accelerating field.
2. It does not require particle bending, which would cause severe synchrotron radiation losses in extreme energies.

3. The accelerating fields and particles move in the collinear direction at the same velocity, the speed of light, so that the acceleration has a built-in coherence called "relativistic coherence" [6]; in contrast, the Fermi acceleration mechanism, based on multiple scatterings, is intrinsically incoherent and stochastic.

4. No escape problem [1] exists. Particles can escape from the acceleration region since the accelerating fields naturally decay out.

5. Whenever and wherever intense electromagnetic waves (with sufficiently high frequencies) are excited, such waves tend to exhibit coherent dynamics (see later for details).

Takahashi et al. [7] and Chen et al. [8] demonstrated that intense Alfvén waves produced by a collision of neutron stars can create wakefields to accelerate charged particles beyond $10^{20}$ eV. Although such a neutron star collision is believed to be related to short gamma-ray bursts [9], it is rather rare for two neutron stars to hit each other directly: It requires the same masses, otherwise the tidal field of the more massive star destroys the less massive one to form an accretion disk. Chang et al. [10] conducted a one-dimensional numerical simulation showing that whistler waves emitted from an AGN produce wakefields to accelerate UHECRs.

The accreting supermassive blackhole, the central engine of an AGN, is one of the candidates for wakefield acceleration. The accretion disk repeats transitions between a highly magnetized (low-beta) state and a weakly magnetized (high-beta) state [11]. In fact, O’Neil et al. [12] have found that magnetic transitions with $10^{-20}$ can lead to the generation of UHECRs beyond $10^{20}$ eV. The paper is organized as follows: We introduce our model for the generation of the intense accelerating structure based on the Tajima-Dawson mechanism [4] that is not hampered by the Fermi mechanism limitations in Section 2, and find that the highest energy is achievable around an accreting black hole (AGN) in Section 3. Astrophysical implications are discussed in Section 4.

2. Intense pondermotive mechanism

An accretion disk is formed around a black hole when gas accretes onto it. Since the angular velocity is higher in inner orbits, there arises a strong shear flow between gases circulating at different radii in the disk. Since the gas is almost fully ionized and Ohmic loss is negligible, magnetic fields are stretched and amplified by the shear motion. The resultant toroidal magnetic field acts as an enhanced friction between gases circulating in the different orbits and transfers the angular momentum outward, while gas is pushed inward because of the reaction of the momentum exchange.

The inner edge of the accretion disk is located around $R = 3R_g$, where

$$R_g = 2GM/c^2 = 3.0 \times 10^{13}(m/10^8)\text{ cm}$$

is the gravitational radius of the blackhole. Here, $m$ is the mass of the blackhole in the unit of solar mass ($2.0 \times 10^{33}$ g). An ergosphere appears just outside of the causality horizon of the blackhole. The gas inside the ergosphere and outside the horizon can extract rotational energy from the blackhole, if it is magnetized. This energy then drives relativistic jets in the two axial directions of the accretion disk [13]. The Lorentz factor $\Gamma$ of the bulk motion of the jet is observed as $10 \sim 30$ in the case of active galactic nuclei.

According to Shibata et al. [11], the accretion disk makes transitions between two states: In the weakly magnetized state, magnetic fields are amplified by a strong shear flow, growing until a certain point, and decay out; in other words, the disk makes transitions between these two states repeatedly. As a result, strong fluctuations are induced in the innermost region of the accretion region ($R < 10R_g$). The physical parameters in this innermost region ($R < 10R_g$) are estimated according to Shukurov and Sunyaev [14]:

$$\dot{\epsilon}_0 = 6.6 \times 10^5(m/10^8)^{-1}\text{ erg cm}^{-3},$$

$$n_0 = 2.9 \times 10^{14}(m/10^8)^{-2}\text{ cm}^{-3},$$

$$Z_0 = 2.2 \times 10^{23}(m/10^8)^{-1}\text{ cm},$$

$$B_0 = 1.8 \times 10^5(m/10^8)^{-1/2}\text{ G},$$

where $m$ is the accretion rate normalized to the critical accretion rate ($\dot{M}_c = L_{\dot{M}0}/0.06c^2$) [14]. The viscosity parameter $\alpha$ is assumed to be 0.1 in the present paper. From the definition of $m$ and $\dot{m}$, the total luminosity of the accreting blackhole is given by

$$L_{\dot{m}} = 1.3 \times 10^{43}(\dot{m}/0.1)(m/10^8)^{-1}\text{ erg s}^{-1}.$$ (6)

The wavelength $\lambda_A$ of Alfvén waves emitted from the accretion disk is calculated as [15]:

$$\lambda_A = \left( V_{ad}/c \right) |\Omega/A| Z_0 = B_0 Z_0 / 3 (4\pi \epsilon_0)^{1/2} = 5.8 \times 10^{12}(m/0.1)(m/10^8)^{-1/2}\text{ cm},$$

where $V_{ad}$ is the Alfvén velocity in the accretion disk, which is calculated as:

$$V_{ad} = B_0 / \sqrt{4\pi m_0 n_0} = 2.4 \times 10^7(m/0.1)$$

and $c$ is the sound velocity in the accretion disk:

$$c = \sqrt{\epsilon_0/m_0 n_0}.$$ (9)
where \( m_0 \) is the proton mass. We assume magnetic field in the accretion disk as \( B_0 \) and the Keplerian rotation of gas inside the disk, i.e. \( \Omega / A = 4/3 \). The magnetic energy \( E_B \) stored in the innermost region of the accretion disk (\( R < 10R_g \)) is estimated as:

\[
E_B = (B_0^2/4\pi)\pi(10R_g)^2Z_D = 1.6 \times 10^{48}(\text{m}/0.1)(\text{m}/10^8)^2 \text{erg.} \quad (10)
\]

The Alfven waves excited in the accretion disk propagate along the global magnetic field of the jet. The normalized vector potential \( a \), which is the Lorentz-invariant strength parameter of the wave [5], is calculated as:

\[
a = eE/m_0c_A, \quad (11)
\]

where \( m_e \) and \( e \) are the electron mass and charge, and we used \( E = (\text{m}_p/\text{c})^{1/2}B_0 \) and \( \omega_m = 2\pi V_A/\lambda \approx 2\pi c/\lambda \). The former comes from the conservation of Alfven energy flux, i.e., \( \Phi_{\text{m}} = eE \times B/4\pi = \Phi_{\text{m}} = V_A B_0 /4\pi \). We find that \( a \) is much greater than unity for a large class of AGN disks. We also find that the Alfven velocity in the jet (except at the very vicinity of the blackhole) is close to \( c \), and thus these Alfven waves exert an intense pondermotive force on electrons and ions. In this \( a \gg 1 \) regime, the longitudinal pondermotive acceleration dominates the transverse acceleration. As we have mentioned in the introduction (page 3), the Tajima-Dawson acceleration [4] requires the conditions (a) and (b). Ashour-Abdalla et al. [16] studied this acceleration mechanism in the astrophysical context, where the condition (b) is overwhelmingly satisfied. It was found [16] that while the pondermotive force accelerates particles ahead of the EM pulse, it causes a density cavity in and behind the pulse (which is the cause of the trailing wakefields). In more recent works with conditions closer to the terrestrial acceleration experiments, Refs. [18–20,17] found qualitatively similar results to Ashour-Abdalla’s, although details vary due to parameter differences. Mourou et al., in his review paper [21], called the EM pulse “relativistic” when \( m_{\text{opt}}/m_0 > a > 1 \), and “ultra-relativistic” when \( a > m_{\text{opt}}/m_0 > 1 \). No terrestrial experiments so far have been performed in the “ultra-relativistic” regime. Only a limited number of theoretical works have been devoted to this regime. Therefore, the details of the dynamics of this regime remain to be investigated in the future. Within 1D, Ashour-Abdalla et al. [16] find that the greater \( a \) is in the “ultra-relativistic” regime, the more that the charge separation force is dominated by the EM pondermotive force, although it is expected that this effect may be mitigated in 2–3D. Thus this regime should be dominated by pondermotive acceleration.

The Alfven flux inside of the jet is assumed to be inversely proportional to \( n \beta^2 \), and \( b \) to the square root of the distance \( D = 10R_g(D/3R_g)^{1/2} \). This scaling is consistent with the VLBI observation of the jet of M87, the closest AGN [22]. In such a case, the value of \( a \) for the wave propagating in the jet is calculated as:

\[
a(D) = a_0(D/3R_g)^{1/2}, \quad (12)
\]

where \( D \) is the distance from the black hole along the jet, and \( a_0 \) is the value of \( a \) at the disk inner edge \( (D = 3R_g) \), which is estimated as:

\[
a_0 = 2.3 \times 10^{30}(\text{m}/0.1)^{3/2}(\text{m}/10^8)^{1/2}. \quad (13)
\]

The Lorentz factor \( \gamma \) of the quivering motion of particles in the wave is of the order of \( a \), i.e., \( \gamma \approx a \). The Alfven pulse generation, its collinear propagation feature, and its pondermotive acceleration all lead to coherent dynamics. In other words, the phase between the specific wave and the particles to be accelerated are tightly locked because the phase velocity of these waves (including the Alfven pulse in the jet under consideration) is very close to the speed of light, and because of the longitudinal (i.e., the direction parallel to the propagation of the Alfven wave which propagates along the direction parallel to the magnetic fields embedded in the jet) nature of the pondermotive force. Further note that the acceleration dynamics in one dimension is robust because of the relativistic coherence [6]. The mechanism known as dephasing (along with the pump depletion) [4,5] determines the maximum energy gain as well as the spectrum [23,8] (see Fig. 2).

We focus on the wave modes propagating parallel to the jet magnetic field, since these modes are effective for the linear acceleration to highest energies. The angular frequency of the Alfven wave is:

\[
\omega_m = 2\pi V_A/\lambda \approx 2\pi c/\lambda \quad (14)
\]

where \( V_A = B_0/\sqrt{4\pi m_e n_0} \) is the Alfven velocity in the jet. If we assume the conservation of magnetic flux in the jet, then the magnetic field \( B_j \) in the jet is scaled as:

\[
B_j = B_0(l_0/3R_g)^{-1}; \quad (15)
\]

the plasma density \( n_j \) in the jet is calculated through the kinetic luminosity \( L_j \) of the jet,

\[
L_j = 4\pi n_j m_e c^3 \beta^2 n_0^2 = 4L_{\text{tot}}; \quad (16)
\]

from which one infers that

\[
\eta_j = 2.6 \times 10^8(\text{m}/0.1)(\text{m}/10^8)^{-1}(\xi/10^{-2}(\Gamma/20)^{-2}(D/3R_g)^{-1} \text{cm}^3. \quad (17)
\]

The effective plasma frequency \( \omega_p \) is calculated as:

\[
\omega_p = (4\pi n_e e^2/m_0^2 \Gamma)^{1/2} \quad (18)
\]

\[
= 2.1 \times 10^{-1}(\Gamma/20)^{-1/2}(\xi/10^{-2})^{1/2}(\text{m}/0.1)^{-1/4}(\text{m}/10^8)^{3/4} \times (D/3R_g)^{-1/4} \text{Hz.} \quad (19)
\]

On the other hand, the effective cyclotron frequency \( \omega_c \) is derived as:

\[
\omega_c = eB_0/m_0c \gamma \quad (20)
\]

\[
= 2.3 \times 10^8(\phi/2.0)(\text{m}/0.1)^{3/2}(\text{m}/10^8)^{-1}(D/3R_g)^{-1/2} \text{Hz.} \quad (21)
\]

As an Alfven wave pulse propagates along the jet, the density and magnetic fields decrease, and accordingly the ratios \( \omega_p/\omega_m \) and \( \omega_c/\omega_m \) plummet, as seen in Fig. 3 (for the case of \( m = 0.1, m = 10^8, \gamma = 20, \) and \( \xi = 10^{-2} \)). As \( \omega_p \) approaches \( \omega_m \), the whistler branch of the Alfven pulse turns into the electromagnetic wave [10] and starts to excite pondermotive and wakefield potentials. The distance \( D_1 \) at which \( \omega_p = \omega_m \) is calculated as:

\[
D_1/3R_g = 1.7 \times 10^7(\Gamma/20)^{-1}(\xi/10^{-2})^2(\text{m}/0.1)^{3}(\text{m}/10^8). \quad (22)
\]

On the other hand, the distance \( D_2 \) at which \( \omega_c = \omega_m \) is calculated as:

\[
D_2/3R_g = 5.1 \times 10^5(\text{m}/0.1)^{1/2}(\phi/2.0)^2, \quad (23)
\]

independent of \( m \) or \( m \). As \( D \) increases, \( \omega_c \) approaches \( \omega_m \). In spite of the cyclotron resonance at \( \omega_c \), most of the wave energy is likely to tunnel from the whistler branch to the upper branch beyond the right-hand cut-off frequency

\[
\omega_c^0 = \left[ (\omega_c^2 + 4\omega_i^2)^{1/2} + \omega_i^2 \right]/2, \quad (24)
\]

which is located above the cyclotron resonance \( \omega_c \) in the case of the cold and linear limit [24]. In addition to the linear evanescent tunneling, the nonlinear nature of the EM waves [16]; \( a \gg 1 \) results in the resonance broadening and the nonlinear tunneling (Fig. 2).
3. Highest energy cosmic rays

The phase velocity of Alfvén wave in the jet is close to the light velocity because of the small $n_J$ compared to $n_D$. In such a case, the particles are accelerated by the pondermotive force parallel to the direction of the wave. The maximum energy $W_{\text{PM}}$ in the observer’s frame of the particles’ gain in the region is calculated as:

$$ W_{\text{max}} = \int_B^{D_3} F_{\text{pm}} dD $$

$$ = 4.6 \times 10^{19} z (\Gamma/20) (m/0.1)^{1/2} (m/10^8)^{1/2} (D_3/3R_g)^{1/2} \text{eV} $$

$$ = 2.9 \times 10^{22} z (\Gamma/20) (m/0.1)^{1/3} (m/10^8)^{2/3} \text{eV}, $$

where

$$ F_{\text{pm}} = \Gamma m c a \omega_{A}, $$

is the pondermotive force of the wave. The acceleration length is assumed to be:

$$ Z_{\text{acc}} = ca / \omega_{A}. $$

This is consistent with Ashour-Abdalla et al. [16]. Further, Barezhani and Murushidze [25] obtained an exact nonlinear longitudinal plasma wave solution excited by a relativistic laser pulse, neglecting the quiver motion of protons. They found that the acceleration length is increased by a factor of $a$ in a fashion similar to Eq. (27). This nature of acceleration lengthening can be expected to remain even in the case that proton quiver motion is not negligible, i.e., $a > 10^3$. They also found the plasma density is significantly reduced in the relativistic laser pulse because the plasma is evacuated by the
strong pondermotive force. Eq. 25 holds as far as $Z_{\text{acc}}$ is greater than $D$. The distance $D_1$ is where the acceleration finishes, defined by the equation

$$D_1 = Z_{\text{acc}} = \frac{AC}{\theta A}.$$  

(28)

We find that particles arrive at $D_1$ before $D_2$, in other words:

$$D_2/3R_\varpi = 3.9 \times 10^{37}(m/0.1)^{0.5}(m/10^9)^{1/3} > D_1/3R_\varpi.$$  

(29)

The energy spectrum of the accelerated charged particles has the power-law with the index of $-2$ in the 1-D model due to the multiple dephasing occurrences when particles ride on and off different peaks of the pondermotive or wakefield hills when the waves contain multiple frequencies (but with again the same phase velocity $\sim c$) [8], i.e., $J(W) = A(W/W_{max})^{-2}$. As noted earlier, when the driving Alfvén waves and their driven pondermotive fields hold a broad band of frequencies, their phase velocities and group velocities, respectively, are again close to the speed of light, providing the basis for the robust accelerating structure. When Alfvén waves have two or three dimensional features, the dephasing is more prompt, leading to higher index of the spectrum (less than $-2$). Let $\kappa$ be the energy conversion efficiency of the acceleration (including the mode convergence efficiency mentioned earlier), then $\kappa E_p = g(W_{max} \ln(W_{max}/W_{min}))$, i.e.,

$$A = 1.6 \times 10^{10} \kappa g \theta A^3 (W_{max} \ln(W_{max}/W_{min}))^{-1}.$$  

(30)

The recurrence rate $v_A$ of the Alfvén pulse burst is evaluated as:

$$v_A = \eta W_{ad}/Z_\varpi = 1.0 \times 10^7 \eta m^{-1} Hz,$$  

(31)

where $\eta$ is episode-dependent, and on the order of unity. This is consistent with the 3-dimensional simulations conducted by O'Neill [12]. They found magnetic fluctuations, called Long Period Quasi Periodic Oscillations (LQPO) with the period 10–20 times the Kepler rotation period. The luminosity $L_{\text{UHECR}}$ of ultra-high energy cosmic rays is:

$$L_{\text{UHECR}} \sim \kappa^2 \varpi v_A = 1.6 \times 10^{37} (\kappa/0.1) \eta m/s^{-1},$$  

(32)

where $\zeta = \ln(W_{max}/10^{37} eV)/\ln(W_{min}/W_{max})$.

The pondermotive fields in the jets accelerate both ions and electrons and therefore the AGN jet is likely to be a strong gamma-ray source as well. Although the radiation loss of protons and nuclei is negligible as far as they are accelerated parallel to the magnetic field [26], that of electrons is likely to be significant, when electrons encounter magnetic fluctuations. The gamma-ray luminosity is, therefore, found to be:

$$L_\gamma \sim \kappa E_p v_A = 1.6 \times 10^{37} (\kappa/0.1) \eta m/s^{-1}.$$  

(33)

We summarize the major features of pondermotive/wakefield acceleration in an accreting supermassive blackhole in Table 1 (Fig. 4).

<table>
<thead>
<tr>
<th>Table 1 Major features of pondermotive acceleration in an accreting supermassive blackhole.</th>
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<tr>
<td><strong>Values</strong></td>
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<tr>
<td>$2\pi/\theta A$</td>
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<tr>
<td>$1/\varpi$</td>
</tr>
<tr>
<td>$D_1/c$</td>
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<tr>
<td>$W_{max}$</td>
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<td>$L_{\text{jet}}$</td>
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<td>$L_{\text{UHECR}}/L_{\text{jet}}$</td>
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<td>$L_{\text{UHECR}}/L_\gamma$</td>
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\[ \zeta = L_{\text{jet}}/L_{\text{UHECR}}. \]

4. Astrophysical implications and blazar characteristics

Radio galaxies belong to one category of AGN, which has radio lobes connected to the nucleus by relativistic jets. Their central engines are accreting supermassive ($m = 10^{4–10^{10}}$) blackholes. Urry and Padovani [27] pointed out that there are parent (or misaligned) populations of blazars, which show rapid time variations in many observational bands across radio to gamma rays (10 GeV) with distinct optical and radio polarizations because of their relativistic jet pointing almost toward us. The recent observation by the Fermi satellite reveals that many blazars emit strong gamma-rays in the GeV energy range [28–30].

We find that radio galaxies are most likely to be sources of UHECRs and their features fit well with the present theory of based on the Tajima-Dawson acceleration. First, according to Ajello et al. [31] and Broderick [32], the local gamma-ray luminosity density of blazars is estimated as $10^{37–38}$ erg s$^{-1}$ (Mpc)$^{-3}$, taking into account the beaming effect of the relativistic jet. Assuming $L_{\text{UHECR}}/L_\gamma \sim \zeta < 1$ (Table 1), our theoretical estimate of UHECR particle flux, averaged over the sky, becomes:

$$\Phi_{\text{UHECR}} = 7.6 \times 10^{-2} L_{\text{jet}}(\zeta/0.1) \times (t_A/1.5) \text{ particles}/(100 \text{ km}^2 \text{ yr sr}).$$  

(34)

Eq. (34) is consistent with observed flux of UHECR. Here, $L_{\text{jet}}$ is the local gamma-ray luminosity density of blazars (in the unit of $10^{37}$ ergs$^{-1}$ (Mpc)$^{-3}$) and $t_A$ is the life time of UHECR particles (in the unit of 10 yr), which is determined by GZK process: Greisen [33] and Zatsepin and Kuzmin [34] predicted that cosmic-ray spectrum has a theoretical upper limit around $5 \times 10^{15}$ eV, because of the opening of the channel to produce $\Delta^+$ particles, which decay into pions ($\pi^+$ and $\pi^-$) and further into photons, electrons, protons, neutrons, and neutrinos. The flux of the cosmicmic neutrinos, produced by the GZK process, is as high as

$$\Phi_{\text{UHEV}} = 5.4 \times 10^{-1} L_{\text{jet}}(\zeta/0.1) \times (t_A/100) \text{ particles}/(100 \text{ km}^2 \text{ yr sr}).$$  

(35)
assuming the conversion efficiency of UHECR to UHE\(\nu\) to be 10\%.
This is consistent with the previous works for the case of \(\nu_{\text{max}} = 10^{21.5}\) \(\text{eV}\) (e.g. [35]). The recently observed PeV neutrinos with IceCube experiment [36] is also consistent if we assume the power law spectrum of the index of \(-2.2\) in the energy region from PeV to ZeV. This level of UHE\(\nu\) flux may be detected by a next generation space borne detector of UHECR, like JEM-EUSO, which can achieve an integrated exposure of 10\(^9\) km\(^2\) str yr [37–40] as well as next generation neutrino facilities in the Antarctic, such as ANITA [41], ARCA [42], and ARIANNA [43].

Second, blazars are also known for being highly variable at all wavelengths and all time scales. In the most extreme cases, the timescales of gamma-ray variability can be as short as a few minutes at very high energies (\(\sim 100\) GeV; VHE). Such variability has been detected in several BL Lacertae objects [36,44,45,29,46,47]. On the other hand, our pondermotive acceleration mechanism predicts the rapid time variability with all the timescales from the Alfven frequency \((2\pi/\nu_{\text{Alfven}} \sim 100\) s\), through the repetition period of the pulses \(1/\nu_{\text{p}} \sim \text{days}\), to the propagation time in the jet \((D_2/c, \sim 1 \sim 10^3 \text{years})\). This time variability is both for ion acceleration variability for UHECRs and even for the electron variability as observed in gamma rays (electron energies are limited by the radiation energy loss by PeV\[48\]). The finer structure of time variability for UHECRs as well as electron variability as observed in gamma rays (energy electrons are limited by the light energy loss by PeV\[48\]). This provides a natural account for UHECRs, and also for accompanying gamma-rays and their related observational characteristics, such as their luminosities, time variations, and structures. The severe physical constraints in the extreme ZeV energies by the Fermi satellite have been lifted by the present mechanism. We have identified a number of areas of future research in need of further studies, including the cavity dynamics of super-intense Alfvenic pulse incurred by an accretion disk as well as by Takahashi et al. [37], Kajino et al. [38], Santangelo [39] and Gorodetzky et al. [40].

5. Conclusions

We have introduced the pondermotive acceleration mechanism arising from the Alfvenic pulse incurred by an accretion disk around a supermassive blackhole, the central engine of an AGN. This provides a natural account for UHECRs, and also for accompanying gamma-rays and their related observational characteristics, such as their luminosities, time variations, and structures. The severe physical constraints in the extreme ZeV energies by the Fermi acceleration have been lifted by the present mechanism. We have identified a number of areas of future research in need of further studies, including the cavity dynamics of super-intense Alfven pulses in \(1-3\) dimensions. We have presented a number of emerging astrophysical phenomena that are not easy to explain by existing theories, but are in line with natural consequences of the present acceleration mechanism.

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