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Time structure and multi-messenger signatures of ultra-high energy cosmic ray sources

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New Journal of Physics **11** (2009) 065014 (13pp)

Received 13 October 2008

Published 30 June 2009

Online at <http://www.njp.org/>

doi:10.1088/1367-2630/11/6/065014

Abstract. The latest results on the sky distribution of ultra-high energy cosmic ray sources have consequences for their nature and time structure, if either deflection is moderate or if their density is comparable to or larger than the average density of active galaxies. If the sources accelerate predominantly nuclei of atomic number A and charge Z and emit continuously, their luminosity in cosmic rays above $\simeq 6 \times 10^{19}$ eV can be no more than a fraction of $\simeq 5 \times 10^{-4} Z^{-2}$ of their total power output. Such sources could produce a diffuse neutrino flux that gives rise to several events per year in neutrino telescopes of km³ size. Continuously emitting sources should be easily visible in photons below ~ 100 GeV, but TeV γ -rays may be absorbed within the source. For episodic sources that accelerate cosmic rays in areas moving with a Lorentz factor Γ , the bursts or flares have to last at least $\simeq 0.1 \Gamma^{-4} A^{-4}$ yr. A considerable fraction of the flare luminosity could then go into highest energy cosmic rays, in which case the rate of flares per source has to be less than $\simeq 5 \times 10^{-3} \Gamma^4 A^4 Z^2 \text{ yr}^{-1}$. Episodic sources should typically have detectable variability both at FERMI/GLAST and TeV energies, but neutrino fluxes may be hard to detect. Finally, in contrast to γ -rays, power and density requirements make it unlikely that the ultra-high energy cosmic rays leave the source environment strongly beamed.

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1. Introduction

Nucleons above $\simeq 70$ EeV suffer heavy energy losses due to photo-pion production on the cosmic microwave background—the Greisen–Zatsepin–Kuzmin (GZK) effect [1]—which limits the distance to possible sources to less than $\simeq 100$ Mpc [2]. A sharp drop in the spectrum, consistent with the GZK effect has indeed been seen recently by both the High Resolution Fly’s Eye (HiRes) [3] experiment and by the Pierre Auger Observatory [4]. The sources of ultra-high energy cosmic rays (UHECR) above this ‘GZK threshold’ are, however, still unknown. Recently, an important step forward has been made by the Pierre Auger Observatory, which has revealed a correlation of the arrival directions of UHECR above $\simeq 6 \times 10^{19}$ eV with the nearby cosmological large scale structure as mapped out by the distribution of active galactic nuclei (AGNs) [5]. It should be noted that such a correlation has not yet been observed in the northern hemisphere [6] which, however, may be due to the considerably smaller exposure.

The amount of UHECR deflection in large scale cosmic magnetic fields is still poorly known. While some studies suggest degree scale deflection angles [7], other studies point out that much larger deflection angles cannot be excluded [8, 9]. If UHECR deflection is moderate, the observed clustering and auto-correlation properties of UHECR events indicate that the source density within the ‘GZK horizon’ of about 75 Mpc is probably higher than a few 10^{-5} Mpc^{-3} . At the same time the observed spectrum normalizes the required injection power per volume. Together, these two numbers imply an upper limit on the time-averaged UHECR injection power per source. Comparing this with the minimal power that needs to be dissipated in order to produce UHECR up to $\sim 10^{20}$ eV, this allows to constrain the time structure of the UHECR emission, in particular, continuously emitting versus episodic sources. This can also have some implications on how the dissipated total power may be distributed between the cosmic ray, photon and neutrino channels. The latter can be important for a multimessenger study of UHECR sources. In the present paper we attempt to work out these constraints in a largely model-independent way, including their dependence on the type of nuclei that are predominantly accelerated. Considering nuclei heavier than nucleons is motivated by the fact that nuclei can be accelerated to higher energies and that the AGNs correlating with the observed UHECR events have been argued to be too weak to accelerate protons to the required energies [10, 11], but may meet the theoretical requirements if accelerated primaries consist of nuclei [12, 13].

The remainder of this paper is structured as follows. In section 2, we develop general requirements on the individual sources. In section 3 and 4, we consider continuously emitting and episodic sources, respectively. In section 5, we discuss Centaurus A as a potential UHECR source and we conclude in section 6. We will use the units in which $c = 1$ throughout.

2. Requirements on individual sources

Accelerating particles of charge eZ to an energy E_{\max} requires an induction $\mathcal{E} \gtrsim E_{\max}/(eZ)$. With $Z_0 \simeq 100 \Omega$ the vacuum impedance, this requires dissipation of a minimal power of [14, 15]

$$L_{\min} \simeq \frac{\mathcal{E}^2}{Z_0} \simeq 10^{45} Z^{-2} \left(\frac{E_{\max}}{10^{20} \text{ eV}} \right)^2 \text{ erg s}^{-1}. \quad (1)$$

We stress that this minimal power can be smaller in specific geometrical circumstances, such as in relativistic blast waves [16]. However, given other, larger uncertainties such as the chemical composition of UHECRs, we will ignore such details in the present work.

The ‘Poynting’ luminosity equation (1) can also be obtained from the expression $L_{\min} \sim \Gamma^2 (BR)^2$, where Γ is the beaming factor of the accelerating region and the product of the size R and magnetic field strength B of the acceleration region is given by the ‘Hillas criterium’ [17] which states that the Larmor radius $r_L = E_{\max}/(\Gamma e Z B)$ should be smaller than R ,

$$\left(\frac{B}{\text{G}} \right) \left(\frac{R}{\text{cm}} \right) \gtrsim 3 \times 10^{17} \Gamma^{-1} \left(\frac{E_{\max}}{Z 10^{20} \text{ eV}} \right). \quad (2)$$

In the following, we denote cosmic rays above $6 \times 10^{19} \text{ eV}$ as UHECR and we take $E_{\max} \simeq 10^{20} \text{ eV}$ as the benchmark for their typical production energy within the sources. Any source producing UHECR up to energy E_{\max} at a given time has to have a total power output of at least the Poynting luminosity equation (1). Note that this is comparable to the Eddington luminosity $L_{\text{Edd}}(M) = 1.3 \times 10^{38} (M/M_{\odot}) \text{ erg s}^{-1}$ of a massive black hole of mass M in the centers of active galaxies. A considerable part L_{γ} of that power is presumably electromagnetic and thus emitted in photons. We now assume that electromagnetic power is produced in the same area of size R in which UHECR are accelerated. Denoting the characteristic photon energy by ε , the optical depth for pion production on such photons by accelerated protons with an energy above the photo-pion threshold, $E \gtrsim 6.8 \times 10^{16} (\varepsilon/\text{eV})^{-1} \text{ eV}$, is given by

$$\begin{aligned} \tau_{p\gamma} &\simeq \sigma_{p\gamma} n_{\gamma} R \simeq \frac{\sigma_{p\gamma} L_{\gamma}}{4\pi R \varepsilon} \\ &\simeq 0.15 \left(\frac{L_{\gamma}}{10^{45} \text{ erg s}^{-1}} \right) \left(\frac{R}{\text{pc}} \right)^{-1} \left(\frac{\varepsilon}{\text{eV}} \right)^{-1}, \end{aligned} \quad (3)$$

where we have used $\sigma_{p\gamma} \simeq 300 \mu\text{barn}$ around the threshold for pion production. Note that $R \sim 1 \text{ pc}$ is the typical size of the accretion disc around supermassive black holes at the centers of AGNs, which is determined by the ‘sphere of influence’ $\sim 2G_N M/v_s^2 \sim 2 (M/10^7 M_{\odot}) (v_s/200 \text{ km s}^{-1})^{-2} \text{ pc}$ where G_N is Newton’s constant and v_s is the velocity dispersion of the stars in the host galaxy (see e.g. [18]). The optical depth for photo-disintegration of primary nuclei is comparable to equation (3). If it is significantly larger than unity, most nuclei would be disintegrated before leaving the source and the maximal energy

would have to be $E_{\max} \gtrsim A 10^{20}$ eV in order for UHECR to arrive at Earth with energies up to 10^{20} eV.

On the other hand, the optical depth for hadronic interactions of accelerated protons and nuclei with the surrounding bulk matter of hadronic mass M_{bulk} extending over a characteristic scale $R_{\text{bulk}} \gtrsim R$ can be written as

$$\begin{aligned} \tau_{\text{pp}} &\simeq \sigma_{\text{pp}} n_{\text{p}} R_{\text{bulk}} \simeq \frac{\sigma_{\text{pp}} M_{\text{bulk}}}{R_{\text{bulk}}^2 m_{\text{N}}} \\ &\simeq 100 \left(\frac{M_{\text{bulk}}}{10^7 M_{\odot}} \right) \left(\frac{R_{\text{bulk}}}{\text{pc}} \right)^{-2}, \end{aligned} \quad (4)$$

where we have estimated the nucleon density by $n_{\text{p}} \sim M_{\text{bulk}}/R_{\text{bulk}}^3$. Note that the mass of AGN accretion discs is roughly comparable with the mass of the central supermassive black hole [18] whose typical mass is $10^{7-8} M_{\odot}$. Equation (4) is only a rough estimate because the details will depend on the geometry, for example spherical versus disc-like accretion. Since the bolometric luminosities of most AGNs are $\ll 10^{47}$ erg s $^{-1}$, a comparison of equations (3) and (4) suggests that hadronic interactions dominate over photo-hadronic interactions if the matter distribution around the cores of AGNs is not strongly clumped.

Pionic and photohadronic processes will produce secondary γ -rays and neutrinos. The optical depth for photons of energy above the pair production threshold, $E \gtrsim m_{\text{e}}^2/\varepsilon \simeq 0.26(\varepsilon/\text{eV})^{-1}$ TeV, can be estimated as

$$\begin{aligned} \tau_{\gamma\gamma} &\simeq \sigma_{\text{T}} n_{\gamma} R \simeq \frac{\sigma_{\text{T}} L_{\gamma}}{4\pi R \varepsilon} \\ &\simeq 300 \left(\frac{L_{\gamma}}{10^{45} \text{ erg s}^{-1}} \right) \left(\frac{R}{\text{pc}} \right)^{-1} \left(\frac{\varepsilon}{\text{eV}} \right)^{-1}, \end{aligned} \quad (5)$$

where $\sigma_{\text{T}} \simeq 0.6$ barn is the Thomson cross section. Note, however, that this optical depth could be strongly reduced if the emission of γ -rays is beamed (see e.g. [19]). In contrast, the charged UHECR may be much less beamed if they are deflected in the environment of the acceleration region. Since magnetic fields of Gauss strength can be present in accretion discs of parsec-scale size [20, 21] and fields on the order of $10 \mu\text{G}$ over kpc scales are common in the host galaxies of AGNs, equation (7) below suggests a significant isotropization of UHECRs in such environments. Below we will, in fact, see that strong UHECR beaming would be hard to reconcile with the number of sources indicated by the latest UHECR observations.

We now deduce some requirements on the size R of the accelerating region. We will also take into account a possible beaming factor Γ such that B and R and other length scales are measured in the comoving frame, whereas luminosities and the energy E_{\max} refer to the observer frame. The synchrotron loss length for a nucleus of atomic number A and charge Z in a magnetic field of strength B is

$$l_{\text{synch}} \simeq 0.43 \Gamma A^4 Z^{-2} \left(\frac{E_{\max}}{10^{20} \text{ eV}} \right)^{-1} \left(\frac{B}{\text{G}} \right)^{-2} \text{ pc}, \quad (6)$$

and the Larmor radius can be written as

$$r_{\text{L}} \simeq 0.1 \Gamma^{-1} Z^{-1} \left(\frac{E_{\max}}{10^{20} \text{ eV}} \right) \left(\frac{B}{\text{G}} \right)^{-1} \text{ pc}. \quad (7)$$

Since any energy loss time must be longer than the acceleration time which itself is larger than the Larmor radius, one has the condition $l_{\text{synch}} \gtrsim r_L$, which gives an upper limit on the magnetic field strength

$$B \lesssim 4.3 \Gamma^2 A^4 Z^{-1} \left(\frac{E_{\text{max}}}{10^{20} \text{ eV}} \right)^{-2} \text{ G}. \quad (8)$$

Together with equation (7) this results in a lower limit on the Larmor radius

$$r_L \gtrsim 2.3 \times 10^{-2} \Gamma^{-3} A^{-4} \left(\frac{E_{\text{max}}}{10^{20} \text{ eV}} \right)^3 \text{ pc}. \quad (9)$$

In case of AGN sources, for example, this is certainly consistent with a size $R \sim 1 \text{ pc}$ for the typical size of an accretion disc. Acceleration could thus occur in small parts of such accretion discs.

It is usually thought that UHECR acceleration happens in the jets of AGNs rather than in the central regions [22, 23] where strong energy loss often prevents acceleration beyond $\sim 10^{18} \text{ eV}$ (see e.g. [24]). The relevant length scales of tens to hundreds of kpc would then imply continuous emission on time scales $\gtrsim 10^5 \text{ yr}$. However, there are also some scenarios for acceleration close to the central supermassive black hole: this can occur, for example, due to the electromotive force induced by magnetic field threading the event horizon where the emitted power is driven by spin-down [25] and the accretion rate could be very small, giving rise to a ‘dead quasar’ [26]. Similarly, UHECR could be accelerated in the polar cap regions of black hole magnetospheres provided the poloidal magnetic fields are slightly misaligned with the black hole spin [27]. More generally, it has been shown that one-shot acceleration of UHECR with energy losses dominated by curvature radiation is possible [12, 28]. These mechanisms are consistent with the estimates discussed above and can potentially lead to UHECR production in flares.

3. Continuously emitting sources

Assuming at most moderate deflection in intergalactic space, the number of arrival directions observed by the Pierre Auger Observatory [29] and other experiments implies a lower limit on the source density [30, 31],

$$n_s \gtrsim n_0 \equiv 3 \times 10^{-5} \Gamma_{\text{CR}}^2 \text{ Mpc}^{-3}, \quad (10)$$

where Γ_{CR} accounts for possible UHECR beaming. This is also consistent with a recent Monte Carlo study of the auto-correlation function of the 27 UHECR events detected by the Pierre Auger experiment [29] above 57 EeV, which deduces a best fit range of $n_s = (4\text{--}14) \times 10^{-5} \text{ Mpc}^{-3}$ [32]. In order to account for the possibility of large deflection, we will in the following not assume the lower limit equation (10), but instead keep the source density explicit.

The UHECR flux observed by the Pierre Auger observatory is [33]

$$\frac{dN_{\text{CR}}}{dE}(E \simeq 6 \times 10^{19} \text{ eV}) \simeq 6 \times 10^{-40} \text{ cm}^{-2} \text{ sr}^{-1} \text{ s}^{-1} \text{ eV}^{-1}, \quad (11)$$

which corresponds to a power per volume of [31]

$$Q_{\text{UHE}} \sim 1.3 \times 10^{37} \text{ erg Mpc}^{-3} \text{ s}^{-1}. \quad (12)$$

Comparing equations (10) and (12) implies for the time-averaged UHECR luminosity per source

$$\overline{L_{\text{UHE}}} \lesssim 4 \times 10^{41} \left(\frac{3 \times 10^{-5} \text{ Mpc}^{-3}}{n_s} \right) \text{ erg s}^{-1}, \quad (13)$$

where $n_s \gtrsim 3 \times 10^{-5} \text{ Mpc}^{-3}$ is valid for small deflection. This is much smaller than the instantaneous total luminosity required by equation (1).

If UHECR sources emit continuously, equations (1) and (13) imply that these sources must emit at least $\simeq 2000 Z^{-2} (n_s/3 \times 10^{-5} \text{ Mpc}^{-3})$ times more energy in channels other than UHECR. This is consistent with the fact that at redshift zero an average AGN in an active state has a bolometric luminosity of $\simeq 5 \times 10^{44} \text{ erg s}^{-1}$, comparable to equation (1), and the volume emissivity is $\simeq 3 \times 10^{40} \text{ erg Mpc}^{-3} \text{ s}^{-1}$, a factor of a few thousand larger than equation (12) [34, 35]. The average AGN bolometric luminosity as inferred from the AGN luminosity function and the AGN volume emissivity correspond to a density of ‘typical’ AGNs of $\simeq 6 \times 10^{-5} \text{ Mpc}^{-3}$, consistent with equation (10). Comparing with equation (10), this implies that strongly beamed UHECR emission is unlikely.

As a result, if sources emit continuously and the total power is distributed roughly equally between hadronic cosmic rays and electromagnetic power, the cosmic ray injection spectrum could extend down to $\lesssim 10^{17} \text{ eV}$ with a rather steep spectrum $\propto E^{-\alpha}$, $\alpha \lesssim 2.7$. Alternatively, if $\alpha \simeq 2.2\text{--}2.3$, the spectrum could reach down to GeV energies. Using equation (12), this can be written as

$$\frac{d\dot{n}_{\text{CR}}}{dE}(E) \simeq (\alpha - 2) \frac{Q_{\text{UHE}}}{(10^{20} \text{ eV})^2} \left(\frac{E}{10^{20} \text{ eV}} \right)^{-\alpha}. \quad (14)$$

Furthermore, equation (4) suggests that the optical depth for hadronic interactions can be of order unity in the cores of AGNs and thus a considerable part of that cosmic ray flux could be transformed to neutrinos with energies $\sim 10^{17} \text{ eV}$. Following [36], for proton primaries, $Z = 1$, we can write for the production rate per volume of neutrinos

$$\frac{d\dot{n}_\nu}{dE}(E) \simeq \frac{2f}{3x_\nu} \frac{d\dot{n}_{\text{CR}}}{dE}(E/x_\nu), \quad (15)$$

where $x_\nu \simeq 0.05$ is the average neutrino energy in units of the parent cosmic ray energy and $f = e^\tau - 1$ is the ratio of number of cosmic rays interacting within the source to cosmic rays leaving the source. If the cosmic ray injection spectrum $\propto E^{-\alpha}$ extends down to E_{min} without break, f is limited by

$$f \left(\frac{E_{\text{min}}}{10^{20} \text{ eV}} \right)^{2-\alpha} \lesssim 2 \times 10^3 Z^{-2} \frac{L_{\text{tot}}}{L_{\text{min}}}, \quad (16)$$

where L_{tot} is the total luminosity and L_{min} is given by equation (1). This condition just results from comparing the total emissivity with the output in UHECR and neutrinos and would be saturated if the total output were dominated by neutrinos in the case of ‘hidden sources’. Since neutrinos do not interact during propagation and ignoring redshift evolution, we can estimate the all-flavor diffuse neutrino flux as

$$j_\nu^{\text{diff}}(E) \simeq \frac{1}{4\pi H_0} \frac{d\dot{n}_\nu}{dE}(E), \quad (17)$$

where $H_0 = 100 h \text{ km s}^{-1} \text{ Mpc}^{-1}$ is the Hubble constant with $h \simeq 0.72$. Putting together equations (14), (15) and (17), we obtain

$$E^2 j_v^{\text{diff}}(E) \simeq 190 x_v^{\alpha-1} (\alpha - 2) f \left(\frac{E}{10^{20} \text{ eV}} \right)^{2-\alpha} \text{ eV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}. \quad (18)$$

If the ankle marks the transition from galactic to extragalactic cosmic rays, the injection spectral index of the latter has to be $\alpha \simeq 2.2$, especially if heavier nuclei are accelerated (see e.g. [37]). The secondary neutrino spectrum would then extend down to at least $\simeq 10^{17} \text{ eV}$ and equation (18) implies

$$E^2 j_v^{\text{diff}}(E) \simeq 4.2 f \left(\frac{E}{10^{17} \text{ eV}} \right)^{-0.2} \text{ eV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}. \quad (19)$$

In contrast, if the ankle is due to pair production of extragalactic protons, then one needs $\alpha \simeq 2.6$ [38]. The secondary neutrino spectrum would then extend down to at least $\simeq 10^{16} \text{ eV}$ and equation (18) implies

$$E^2 j_v^{\text{diff}}(E) \simeq 240 f \left(\frac{E}{10^{16} \text{ eV}} \right)^{-0.6} \text{ eV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}. \quad (20)$$

We now compute the neutrino event rates in kilometer scale neutrino observatories for these two scenarios. Using the neutrino–nucleon cross section $\sigma_{\nu N} \simeq 1.9 \times 10^{-33} (E/10^{16} \text{ eV})^{0.363} \text{ cm}^2$ for $10^{16} \text{ eV} \lesssim E \lesssim 10^{21} \text{ eV}$ [39], we obtain the rate

$$\begin{aligned} R_v &\sim \sigma_{\nu N}(E) 2\pi E j_v^{\text{diff}}(E) n_N V_{\text{eff}} \\ &\sim 2.3 \left(\frac{E}{10^{16} \text{ eV}} \right)^{-0.637} \left(\frac{V_{\text{eff}}}{\text{km}^3} \right) \left(\frac{E^2 j_v^{\text{diff}}(E)}{100 \text{ eV cm}^{-2} \text{ sr}^{-1} \text{ s}^{-1}} \right) \text{ yr}^{-1}, \end{aligned} \quad (21)$$

where $n_N \simeq 6 \times 10^{23} \text{ cm}^{-3}$ is the nucleon density in water/ice and V_{eff} is the effective detection volume. The scenario $\alpha = 2.2$ with neutrino flux down to 10^{17} eV would give $\simeq 2 \times 10^{-2} f \text{ yr}^{-1} \text{ km}^{-3} \lesssim 20 Z^{-2} (L_{\text{tot}}/L_{\text{min}}) \text{ yr}^{-1} \text{ km}^{-3}$, where we have used equation (16) for f . The scenario $\alpha = 2.6$ with neutrino flux down to 10^{16} eV would give $\simeq 5.5 f \text{ yr}^{-1} \text{ km}^{-3} \lesssim 180 Z^{-2} (L_{\text{tot}}/L_{\text{min}}) \text{ yr}^{-1} \text{ km}^{-3}$.

According to equation (5), TeV γ -rays would be visible only for considerably beamed emission. As opposed to [36], we therefore do not necessarily get a constraint from the non-observation of AGNs at TeV energies in this scenario. In contrast, x-rays and GeV γ -rays could leave the source. Individual sources should be visible by x-ray telescopes and by FERMI/GLAST. In fact, EGRET has seen a diffuse flux [40], which constrains the neutrino flux because a comparable amount of energy goes into photons and neutrinos in primary cosmic ray interactions,

$$E^2 j_v^{\text{diff}}(E) \lesssim 10^3 \text{ eV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}. \quad (22)$$

Equations (19) and (20) satisfy this limit except for sources deeply in the hidden regime, $f \gg 1$.

4. Flaring sources

If the sources flare on a typical time scale δt in the observer frame, the corresponding life time of the burst in the comoving frame, $\Gamma \delta t$ must be larger than the comoving acceleration time scale, which itself is larger than the Larmor radius, thus with equation (9) we have

$$\delta t \gtrsim 0.1 \Gamma^{-4} A^{-4} \left(\frac{E_{\max}}{10^{20} \text{ eV}} \right)^3 \text{ yr}. \quad (23)$$

This is consistent with the variabilities observed for AGNs, which are observed at time scales down to ~ 60 s [41], provided that $\Gamma \gtrsim 20$ and/or predominantly heavier nuclei are accelerated. The time scale is also consistent with γ -ray bursts, which can easily have Lorentz factors $\Gamma \gtrsim 20$ [42].

If we denote the rate and UHECR luminosity of typical flares by R_f and L_{UHE} , respectively, we can write for the time-averaged UHECR power, equation (13),

$$R_f \delta t L_{\text{UHE}} \sim \overline{L_{\text{UHE}}} \lesssim 4 \times 10^{41} \left(\frac{3 \times 10^{-5} \text{ Mpc}^{-3}}{n_s} \right) \text{ erg s}^{-1}. \quad (24)$$

This means that the fraction of time a typical intermittent source emits the typical instantaneous UHECR luminosity L_{UHE} , also called the *duty factor*, is given by

$$\mathcal{D} \equiv R_f \delta t \lesssim 4 \times 10^{-4} Z^2 \left(\frac{3 \times 10^{-5} \text{ Mpc}^{-3}}{n_s} \right) \left(\frac{L_{\min}}{L_{\text{UHE}}} \right) \left(\frac{E_{\max}}{10^{20} \text{ eV}} \right)^{-2}. \quad (25)$$

Equations (24) and (25) imply that a considerable fraction of the minimal flaring luminosity equation (1) can go into UHECR, $L_{\text{UHE}} \sim L_{\min}$, which could be the case if the UHECR acceleration spectrum is hard, $\alpha \lesssim 2$. From equations (24) and (23) we then have

$$R_f \lesssim 5 \times 10^{-3} \Gamma^4 A^4 Z^2 \left(\frac{3 \times 10^{-5} \text{ Mpc}^{-3}}{n_s} \right) \left(\frac{L_{\min}}{L_{\text{UHE}}} \right) \left(\frac{E_{\max}}{10^{20} \text{ eV}} \right)^{-3} \text{ yr}^{-1}. \quad (26)$$

In the limit $L_{\text{UHE}} \rightarrow \overline{L_{\text{UHE}}}$ we obviously recover the limit of continuous sources, $R_f \delta t \rightarrow 1$.

During one flare the total non-thermal energy release would be

$$E_f \gtrsim L_{\min} \delta t \gtrsim 3 \times 10^{51} \Gamma^{-4} A^{-4} Z^{-2} \left(\frac{E_{\max}}{10^{20} \text{ eV}} \right)^5 \text{ erg}, \quad (27)$$

which is consistent with the estimates in [43]. If this energy release is due to accretion onto a central black hole with an energy extraction efficiency of $\sim 10\%$ [44], this corresponds to about $0.02 \Gamma^{-4} A^{-4} Z^{-2} M_{\odot}$. In AGN scenarios involving accretion, this energy requirement is certainly modest.

Individual sources would be observed with the apparent UHECR luminosity

$$L_{\text{UHE,obs}} \simeq L_{\text{UHE}} \frac{\delta t}{t_{\text{disp}}}. \quad (28)$$

Here, t_{disp} is the time dispersion of charged cosmic rays due to deflection in cosmic magnetic fields which for propagation over distance d in a stochastic extragalactic magnetic field $B_{\text{eg}} \sim \text{nG}$ of coherence length $l_c \sim 1 \text{ Mpc}$ can be estimated as

$$t_{\text{disp}} \simeq 4 \times 10^3 Z^2 \left(\frac{E}{6 \times 10^{19} \text{ eV}} \right)^{-2} \left(\frac{d}{10 \text{ Mpc}} \right)^2 \left(\frac{l_c}{1 \text{ Mpc}} \right) \left(\frac{B_{\text{eg}}}{\text{nG}} \right) \text{ yr}. \quad (29)$$

The corresponding deflection angle

$$\theta \simeq 1.3^\circ Z \left(\frac{E}{6 \times 10^{19} \text{ eV}} \right)^{-1} \left(\frac{d}{10 \text{ Mpc}} \right)^{1/2} \left(\frac{l_c}{\text{Mpc}} \right)^{1/2} \left(\frac{B_{\text{eg}}}{\text{nG}} \right) \quad (30)$$

would actually be sufficiently small for the lower limit on the source density equation (10) to apply even for a heavy UHECR composition [32]. Note that for magnetic fields with homogeneous statistical properties, current upper limits from Faraday rotation measurements are $B \lesssim 2.8 \times 10^{-7} (\Omega_b h^2 / 0.02)^{-1} (h / 0.7) (l_c / \text{Mpc})^{-1/2} \text{ G}$, where the baryon density $\Omega_b h^2 \simeq 0.02$ and h is the Hubble constant in $100 \text{ km s}^{-1} \text{ Mpc}^{-1}$ [45].

TeV γ -rays may or may not be observable from individual sources because the duty cycle is small and most of the time the source would have luminosities $\ll 10^{45} Z^{-2} \text{ erg s}^{-1}$, leading to fluxes $\ll 2 \times 10^{-8} Z^{-2} (d / 20 \text{ Mpc})^{-2} \text{ erg cm}^{-2} \text{ s}^{-1}$, where d is the distance to the source. In the active phases, equation (5) suggests that TeV γ -ray may be absorbed by pair production within the sources, but x-rays would not suffer much absorption. Note that flares in the electromagnetic luminosity are not expected to correlate with the UHECR luminosities due to the large UHECR time delays.

For a number of UHECR flares per volume and time \dot{n}_{UHE} equation (12) implies

$$\dot{n}_{\text{UHE}} \delta t \simeq 1.3 \times 10^{-8} \left(\frac{L_{\text{UHE}}}{10^{45} \text{ erg s}^{-1}} \right)^{-1} \text{ Mpc}^{-3}. \quad (31)$$

A recent study compared this required rate with luminosity functions and upper limits on bright extragalactic objects in x- and γ -rays [46]. They concluded that as long as the emission power going into the electromagnetic channel is not considerably smaller than L_{UHE} , the transient flare power should satisfy $L_\gamma \gtrsim 10^{50} \text{ erg s}^{-1}$.

In the flaring limit, we can have $L_{\text{UHE}} \sim L_{\text{min}}$ in which case both γ -ray fluxes and secondary neutrino fluxes cannot be much larger than the UHECR flux, which would then be a considerable fraction of the total energy budget. Since the spectrum must be rather hard in this case, both the diffuse and discrete neutrino fluxes are likely unobservably small.

5. Centaurus A and other AGN sources

Centaurus A is the nearest AGN of the Fanaroff–Riley type I, at a distance $d \simeq 4 \text{ Mpc}$ with a central supermassive black hole of mass $M \sim 10^8 M_\odot$ [47]. It is not a blazar as its jet has a large inclination angle to the line of sight. The Pierre Auger Observatory measured two UHECR events from the direction of Centaurus A [5, 29] and several more events are in fact aligned with its giant radio lobes [10, 48]. This corresponds to a flux [49]

$$\frac{dN_{\text{CR}}}{dE}(E \simeq 6 \times 10^{19} \text{ eV}) \lesssim 10^{-40} \text{ cm}^{-2} \text{ sr}^{-1} \text{ s}^{-1} \text{ eV}^{-1}, \quad (32)$$

and to an apparent UHECR luminosity of Centaurus A of $L_{\text{UHE,obs}} \lesssim 10^{39} \text{ erg s}^{-1}$. If Cen A is a continuous UHECR source, this is consistent with equation (13). The bolometric luminosity of Cen A is $L_{\text{tot}} \simeq 10^{44} \text{ erg s}^{-1}$, which originates mostly within $\simeq 500 \text{ pc}$ from the center and is mostly emitted around energies $\varepsilon \sim 1 \text{ eV}$ [47, 50]. This is consistent with equation (1) if predominantly heavier nuclei are accelerated, $Z \gtrsim 4$. At MeV energies Cen A has a luminosity $\simeq 10^{42} \text{ erg s}^{-1}$ [47].

It has been argued that Centaurus A is capable of accelerating UHECR in its giant lobes [51], at least if it has been more active in the past [48], and may also be detectable in

GeV–TeV γ -rays by instruments such as the recently launched FERMI/GLAST observatory [51] and also in ground-based γ -ray telescopes [52].

If Cen A is emitting UHECR continuously, and assuming the UHECR injection spectrum is $\propto E^{-\alpha}$, the expression analog to equation (15) for a discrete source gives for the secondary neutrino flux

$$j_\nu(E) \simeq \frac{2f}{3x_\nu} \frac{dN_{\text{CR}}}{dE}(E/x_\nu). \quad (33)$$

Using equation (32), one obtains numerically for the neutrino energy flux

$$E^2 j_\nu(E) \simeq 0.24 x_\nu^{\alpha-1} f \left(\frac{E}{6 \times 10^{19} \text{ eV}} \right)^{2-\alpha} \text{ eV cm}^{-2} \text{ s}^{-1}. \quad (34)$$

For $\alpha \simeq 2.2$ and $x_\nu \simeq 0.05$, this gives

$$E^2 j_\nu(E) \simeq 0.024 f \left(\frac{E}{10^{17} \text{ eV}} \right)^{-0.2} \text{ eV cm}^{-2} \text{ s}^{-1}, \quad (35)$$

where from equation (16), $f \lesssim 100$. For $\alpha \simeq 2.6$, it yields

$$E^2 j_\nu(E) \simeq 0.37 f \left(\frac{E}{10^{16} \text{ eV}} \right)^{-0.6} \text{ eV cm}^{-2} \text{ s}^{-1}, \quad (36)$$

where from equation (16), $f \lesssim 3.2$. The latter, more optimistic case would give an event rate of $\simeq 1.3 \times 10^{-3} f \text{ yr}^{-1} \text{ km}^{-3} \lesssim 4.2 \times 10^{-3} \text{ yr}^{-1} \text{ km}^{-3}$. It is clear that the rate due to the diffuse flux, estimated below equation (21), is always much larger as long as Cen A is an ‘average’ source. This is consistent with the conclusion in [53].

If Cen A is an episodic UHECR source, as may be suggested by its variability on time scales of days observed in x- and γ -rays [47], equations (28) and (23) imply for the UHECR luminosity during a flare

$$L_{\text{UHE}} \lesssim 10^{43} \frac{\Gamma^4 A^4 Z^2}{\Gamma_{\text{CR}}^2} \left(\frac{t_{\text{disp}}}{Z^2 10^3 \text{ yr}} \right) \left(\frac{E_{\text{max}}}{10^{20} \text{ eV}} \right)^{-3} \text{ erg s}^{-1}. \quad (37)$$

Comparing with equation (1) this suggests that the UHECR flare luminosity is comparable with the total output, as long as the flare duration is not much larger than the theoretical minimal variability time scale equation (23) and typical time dispersion due to large scale magnetic fields is $\gtrsim 10^3 \text{ yr}$.

The closest blazars whose jets are close to the line of sight and thus may have considerably beamed emission are in general too far away to be responsible for UHECR. As an example, we briefly discuss Markarian 501. This blazar at a distance $d \simeq 130 \text{ Mpc}$ shows emission up to TeV energies with variability on time scales of days and peak luminosities of close to $10^{46} \text{ erg s}^{-1}$ [54, 55]. This is consistent with equation (1) even for protons, $Z = 1$. The power of such blazars is, therefore, certainly sufficient to provide the UHECRs. Even if they produce UHECR only in flares, the flare luminosity in UHECR, L_{UHE} , could be a small fraction of the total flare luminosity, and the necessary flaring time scale equation (23) and rate equation (26) would be consistent with observations, especially for significant γ -ray beaming factors Γ typical for blazars.

6. Conclusions

We have discussed some consequences of the latest results on UHECR for the nature and variability of the sources as well as for the secondary γ -ray and neutrino fluxes produced within the sources. To this end we assumed predominant acceleration of nuclei of atomic mass A and charge Z . On the one hand, if deflection in cosmic magnetic fields is moderate, the observed clustering and auto-correlation properties of arrival directions by experiments such as the Pierre Auger Observatory suggests source densities $n_s \gtrsim 3 \times 10^{-5} \Gamma_{\text{CR}}^2 \text{Mpc}^{-3}$ with Γ_{CR} a possible beaming factor of cosmic rays. On the other hand, the total power that needs to be dissipated to produce the highest energy cosmic rays makes much higher densities of suitable sources unlikely. This also implies that UHECR emission is unlikely to be strongly beamed.

In the limit of continuously emitting sources, their luminosity in cosmic rays above $\simeq 6 \times 10^{19} \text{eV}$ can be no more than a fraction of $\simeq 5 \times 10^{-4} Z^{-2} (3 \times 10^{-5} \text{Mpc}^{-3} n_s^{-1})$ of the total source power. If these cosmic rays are produced in the accretion disks in the centers of AGNs, significant neutrino fluxes could be produced by hadronic interactions, especially in scenarios in which extragalactic protons dominate down to $\simeq 10^{17} \text{eV}$ such that the ankle is due to pair production of these protons. The resulting cosmological diffuse neutrino flux can lead to detection rates up to several events per year and km^3 of effective detection volume. This also implies considerable photon fluxes at energies up to $\sim 100 \text{GeV}$, the latter of which should be easily visible by FERMI/GLAST, whereas TeV γ -rays may be absorbed within the source. For episodic sources that are beamed by a Lorentz factor Γ , individual flares have to last at least $\simeq 0.1 \Gamma^{-4} A^{-4} \text{yr}$. Such flares can also be visible in photons up to the TeV energy range. A considerable fraction of the flare luminosity could go into highest energy cosmic rays, which suggests a hard injection spectrum. In this case, the rate of flares per source has to be $\lesssim 5 \times 10^{-3} \Gamma^4 A^4 Z^2 (3 \times 10^{-5} \text{Mpc}^{-3} n_s^{-1}) \text{yr}^{-1}$. In contrast to continuously emitting sources, both neutrino fluxes from individual sources and the resulting cosmological diffuse flux may be hard to detect in the limit of flaring sources. Conversely, if high energy neutrinos are detected in the near future, this may suggest sources that produce UHECR continuously.

Acknowledgments

I thank Peter Biermann, Michael Kachelriess, Jim Matthews, Rachid Ouyed, Georg Raffelt, Dmitry Semikoz and Leo Stodolsky for useful discussions. I acknowledge support by the DFG (Germany) under grant SFB-676. I thank the Max Planck Institut für Physik in Munich for financial support of a visit during which part of this paper was written.

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