Radio Quiet AGNs as Possible Sources of Ultra-high Energy Cosmic Rays

Asaf Pe’er\textsuperscript{1,2}, Kohta Murase\textsuperscript{3} and Peter Mészáros\textsuperscript{4}

\textsuperscript{1}Space Telescope Science Institute, Baltimore, MD 21218, USA
\textsuperscript{2}Riccardo Giacconi Fellow
\textsuperscript{3}Yukawa Institute for Theoretical Physics, Kyoto University, Kyoto, 606-8502, Japan
\textsuperscript{4}Department of Astronomy \& Astrophysics; Department of Physics; Center for Particle Astrophysics; Pennsylvania State University, University Park, PA 16802, USA

Active galactic nuclei (AGNs) have been one of the most widely discussed sources of ultrahigh-energy cosmic rays (UHECRs). The recent results of Pierre Auger observatory (PAO) have indicated a possible composition change of UHECRs above $\sim 10^{18.5}$ eV towards heavy nuclei. We show here that if indeed UHECRs are largely heavy nuclei, then nearby radio quiet AGNs can also be viable sources of UHECRs. We derive constraints on the acceleration sites which enable acceleration of UHECRs to $10^{20}$ eV without suffering losses. We show that the acceleration of UHECRs and the survival of energetic heavy nuclei are possible in the parsec scale weak jets that are typically observed in these objects, the main energy loss channel being photodisintegration. On this scale, energy dissipation by shock waves resulting from interactions inside a jet or of the jet with surrounding material are expected, which may accelerate the particles up to very high energies. We discuss the possible contribution of radio-quiet AGNs to the observed UHECR flux, and show that the required energy production rate in UHECRs by a single object could be as low as $\approx 3 \times 10^{39} \text{erg s}^{-1}$, which is less than a percent of the bolometric luminosity, and thus energetically consistent. We discuss consequences of this model, the main one being the difficulty in detecting energetic secondaries ($\gamma$-rays and neutrinos) from the same sources.

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I. INTRODUCTION

The origin of ultrahigh-energy cosmic rays (UHECRs), cosmic rays with energies above $\sim 10^{18.5}$ eV, is still a mystery. It is commonly believed that the sources of UHECRs are extragalactic, since at these energies the cosmic rays cannot be confined by the magnetic field of our galaxy (for recent reviews, see Refs. \textsuperscript{11,12}). Due to energy losses by photo-meson production with the cosmic microwave background (CMB), the sources of UHECR’s are further limited to a distance $\lesssim 100$ Mpc from the observer (this limit is known as the ‘Greisen-Zatsepin-Kuzmin (GZK) cutoff’; Refs. \textsuperscript{3,4}). While early observations \textsuperscript{5} suggested a violation of the GZK cutoff at high energies, recent, high statistics observations by the High-Resolution Fly’s eye (HiRes) experiment and the Pierre Auger observatory (PAO) finds a downturn in the spectrum consistent with the GZK predictions \textsuperscript{6,7}.

Extracting information from observations is, however, not easy. For example, observations of the spectrum itself did not lead, so far, to a firm conclusion about the origin of the steepening at high energies (i.e., whether it originates from the GZK cutoff or not), and therefore other possibilities are still allowed.

One of the consequences of this difficulty is the lack of a clear theoretical picture of the extragalactic UHECR sources. Possible candidates of UHECR sources are limited among the known astrophysical objects, because of the difficulty in satisfying the two nearly contradictory requirements: A strong magnetic field in the source is needed to confine the accelerated cosmic rays, while the magnetic and photon fields cannot be too strong in order to avoid too much synchrotron radiation and photo-meson energy losses \textsuperscript{8}. Several possible sources of UHECRs that fulfill the constraints mentioned above are often discussed in the literature. The most widely discussed candidates are gamma ray bursts (GRBs) \textsuperscript{9,10,11,12,13,14}, low luminosity GRBs and hypernovae \textsuperscript{14,15,16}, and jets in radio-loud (RL) AGNs \textsuperscript{17,18,19,20,21,22,23}. UHECR acceleration in the vicinity of black holes has also been considered for AGNs, including radio-quiet (RQ) AGNs (see, e.g., Refs. \textsuperscript{24,25}). Additional suggestions include magnetars \textsuperscript{26,27} and clusters of galaxies \textsuperscript{28,29}.

An important step towards understanding the origin of UHECRs was taken recently with the discovery of a correlation between the arrival direction of the highest energy cosmic rays and nearby galaxies in the \textsuperscript{30} (VCV) catalog of active galactic nuclei (AGNs) [12th edition] \textsuperscript{31,32}. These results, however, should be taken with great care. It was noted already by Refs. \textsuperscript{31,32} that the observed correlation alone is insufficient to conclude that the sources of UHECRs are AGNs, because AGNs themselves are clustered within galaxies, and thus may only be the tracers of the true sources. Indeed, a correlation was found between the arrival directions of UHECRs and galaxies \textsuperscript{33,34,35}, which suggests a correlation with the large scale structure (LSS) in the nearby universe.

The emerging picture gets further complicated by the recent results of HiRes experiment which do not confirm the correlation between the arrival directions of the highest energy cosmic rays and nearby galaxies in the \textsuperscript{30} (VCV) catalog of active galactic nuclei (AGNs) [12th edition] \textsuperscript{31,32}. These results, however, should be taken with great care. It was noted already by Refs. \textsuperscript{31,32} that the observed correlation alone is insufficient to conclude that the sources of UHECRs are AGNs, because AGNs themselves are clustered within galaxies, and thus may only be the tracers of the true sources. Indeed, a correlation was found between the arrival directions of UHECRs and galaxies \textsuperscript{33,34,35}, which suggests a correlation with the large scale structure (LSS) in the nearby universe.
the correlation found by PAO. Moreover, the latest PAO results suggest that the correlation degree is reduced compared to that obtained by earlier measurements.

An additional clue on the nature of UHECR sources may come from measurements of their chemical composition. However, deducing the chemical composition of UHECRs from current observations is also difficult due to our poor knowledge on hadronic interactions. PAO results indicate a heavy composition at the highest energies, while HiRes results suggest a proton composition (see Ref. 2, and references therein). The results of these two experiments are therefore clearly in contradiction. Moreover, the heavy element composition suggested by PAO results may be in contradiction with the anisotropy result of the same experiment (see below).

The chemical composition of UHECRs may be crucial in determining the nature of their source. If UHECRs are protons, the theoretical restrictions on the physical conditions inside potential sources are rather extreme, so that if indeed AGNs are the sources of UHECRs, then only powerful AGNs such as RL AGNs can be viable sources. In addition, the known values of the galactic magnetic field, as well as observational upper limits on the intergalactic magnetic fields (IGMFs) in voids 10 and theoretical investigations about IGMFs in the structured region 11 allow one to expect a correlation between the arrival directions of UHECRs and AGNs 12, which is tentatively suggested by PAO. On the other hand, a drawback of the idea that AGNs are sources of UHECRs is the fact that a strong correlation with RL AGNs is not found, except for very few cases, such as Cen A or Cen B 13. In addition, a correlation with the most powerful AGNs, FR II galaxies, is currently not confirmed. Moreover, the number density of FR II galaxies seems too small to explain the observed small scale anisotropy 14. Furthermore, the observed bolometric luminosity of AGNs which are correlated with the arrival direction of UHECRs seems too small to satisfy the minimum conditions for UHECRs acceleration in continuous jets 15. These results indicate that only very few nearby AGNs are sufficiently powerful to accelerate particles to the observed ultrahigh energies.

In order to overcome this problem, it was suggested by Ref. 16 that UHECR production may occur transiently during giant AGN flares caused by tidal disruptions. Although this may be a valid scenario, several restrictions on the duration and rate of such flares can be derived 17, 18. Since the actual activity of AGNs is not known, concrete conclusions cannot be drawn at this stage (see also Ref. 19).

An alternative picture emerges if UHECRs are composed of heavier elements, such as iron nuclei. The ambiguity in the interpretation of the data led 20 to study a mixed composition scenario above $\sim 10^{18.5}$ eV. A similar model (albeit with somewhat lower maximum energy of UHECR) was studied recently by 21. In these models, significant correlation between the arrival directions of UHECRs and galaxies is not expected, because of the strong deflection of heavy nuclei in galactic magnetic field. In addition, protons dominate the cosmic rays spectra only up to energies $\sim 10^{18.5}$ eV, which implies that the requirement on the physical conditions at the acceleration site of potential sources is significantly loosed.

As we show here, a heavy element composition of UHECRs opens the possibility for RQ AGNs, which compose the majority of the AGN population within $\sim 100$ Mpc, to be the sources of UHECRs. As opposed to RL AGNs, in which evidence exist for strong jets on $\gtrsim$ kpc scale, in RQ AGNs there are no evidence of such strong jets 92. Nonetheless, in recent years there has been an accumulation of evidence for weak, parsec-scale jets in these objects (Refs. 54, 55, 56, 57, 58, 59, and references therein). Internal interactions of blobs inside a jet, or of the jets with their environment, produce shock waves, which can accelerate particles. Alternatively, cosmic rays may be accelerated by the dissipation of magnetic energy inside these jets (see, e.g., Ref. 60). Moreover, as was pointed out by Ref. 61, particle acceleration could take place in shocks within the corona, which could be caused by abortive jets.

In this work we thus discuss weak jets in RQ AGNs as possible sources of UHECRs, under the assumption that at very high energies the composition of UHECRs is dominated by heavy (possibly, but not necessarily, iron) nuclei. As we will show below, RQ AGNs can fulfill the requirements needed from sources of UHECRs under this assumption. We assume here that particles acceleration takes place inside the jets, so that this work is different from previous works that considered UHECR production in the vicinity of black holes that do not have jets (e.g., Ref. 22).

This paper is organized as follows. In II we present the basic constraints on the physical conditions at the acceleration sites of heavy nuclei. In particular, in II A we derive the constraints from photodisintegration and photomeson production. We show that these conditions can be fulfilled easily. In II B we discuss the efficiency of particle acceleration required to explain the observed flux of UHECRs on earth, under the assumption that RQ AGNs are the only sources of UHECRs. We further derive constraints on the value of the magnetic field at the acceleration site, and show that a lower limit of few percents of equipartition value is required. We summarize and conclude in IV.

II. CONSTRAINTS ON THE EMISSION SITE OF UHECRS

The acceleration sites of UHECRs are required to fulfill two basic conditions (see, e.g., Refs. 8, 10, 12). The first condition is that the accelerated particles should be confined to the acceleration region (this is also known as the “Hillas condition”, see Ref. 8). Assuming that the ac-
acceleration results from electromagnetic processes within an expanding plasma, this condition is equivalent to the requirement that the acceleration time is shorter than the dynamical time scale. Considering jetted plasma, which moves relativistically with velocity $\beta c$, or Lorentz factor $\Gamma = (1 - \beta^2)^{-1/2}$, the acceleration time scale in the co-moving frame can be written as $t_{\text{acc}} \simeq \eta E^b / (\Gamma Z q B c)$ \cite{53}. Here, $E^b = \Gamma E'$ is the observed energy of the particle \cite{54}, $E'$ is its energy in the co-moving frame, $B$ is the magnetic field, and $Zq$ is the nuclear charge. The factor $\eta \geq 1$ is a dimensionless factor, whose exact value is determined by the (yet uncertain) details of the acceleration mechanism. For example, in the non-relativistic diffusive shock acceleration mechanism, this factor corresponds to $\eta = (20/3) \beta^2$ in the Bohm limit for parallel shocks \cite{16, 62, 63}. For relativistic shocks, $\eta \sim$ a few can be expected but larger values are also possible. The condition $t_{\text{acc}} < t_{\text{dyn}} = r / \Gamma \beta c$, gives

$$B \geq \frac{\eta E^b \beta}{Z q r} = \frac{3.3 \times 10^{17} \eta \beta}{r} Z^{-1} E^b_{20} G, \quad (1)$$

where $r$ is the radial distance from the source at which particle acceleration takes place, measured in cm. Here and below, we use the convention $Q = 10^{5} Q_X$ in cgs units.

Writing the energy density in the magnetic field as a fraction $\epsilon_B$ of the photon energy density at the acceleration site $u_B \equiv B^2 / 8 \pi = \epsilon_B u$, where $u = L / (4 \pi r^2 \Gamma^2 \beta c)$ is the photon energy density (representative of the electron energy density) and $L$ is the photon luminosity, equation \textbf{1} can be written in the form (Ref. \textbf{10})

$$\epsilon_B L \geq 1.7 \times 10^{45} \eta^2 Z^{-2} \Gamma^2 \beta^3 E^b_{20} \text{ erg s}^{-1}. \quad (2)$$

Equations \textbf{1} \textbf{2} give basic conditions required from a source capable of accelerating particles to the maximum observed UHECR energy, $\sim 10^{20}$ eV. However, as the details of the acceleration process are uncertain, the actual restriction may depend on additional phenomena, such as the details of particle escape from the acceleration region. The details of particle escape are unknown, and the maximum energy may be escape-limited if the escape time scale is shorter than the dynamical time scale. In this paper, we simply assume that acceleration is efficient, such that $t_{\text{esc}} \gtrsim t_{\text{dyn}}$, which can be expected when the Bohm limit is achieved over the size of $r / \Gamma$.

For protons ($Z=1$), equation \textbf{4} necessitates a magnetic luminosity which is 1-2 orders of magnitude above the bolometric photon luminosity observed in nearby AGNs associated with UHECR’s in the VCV catalog, $L_{\text{bol}}^b \approx 10^{44} \text{ erg s}^{-1}$ \cite{15}. We therefore expect that for proton composition of UHECR, acceleration to the highest energies is possible only in powerful RL AGNs (for which $L \gtrsim 10^{44} \text{ erg s}^{-1}$). However, for heavier nuclei, equation \textbf{2} implies a much weaker constraint: e.g., for Carbon nuclei ($Z = 6$) one obtains $\epsilon_B L \gtrsim 5 \times 10^{43} \eta^2 \Gamma^2 \beta^3 E^b_{20} \text{ erg s}^{-1}$. For iron nuclei, $Z = 26$ and the constraint on the luminosity further eases to $\epsilon_B L \gtrsim 2.5 \times 10^{42} \eta^2 \Gamma^2 \beta^3 E^b_{20} \text{ erg s}^{-1}$, well within the limits of many nearby AGNs’ observed luminosity (except for low-luminosity AGNs, with $L \lesssim 10^{42} \text{ erg s}^{-1}$) \cite{55}.

The second condition is that the acceleration time is shorter than all the relevant energy loss time scales such as synchrotron loss time. For the synchrotron loss time, we have $t_{\text{cool, syn}} = (6 \pi m_e^3 c^3 / \sigma_T m_e^3 B^2 E^b Z^4)$ (see, e.g., Ref. \textbf{65}). Here, $\sigma_T$ is Thomson’s cross section, $m_p$ and $m_e$ are the proton and electron masses and $A m_p$ is the mass of the nucleon (for iron nuclei, $A = 56$ and $Z = 26$). The requirement $t_{\text{acc}} < t_{\text{cool, syn}}$ results in

$$E^b \lesssim 2 \times 10^{20} \Gamma \eta^{-1/2} B^{-1/2} Z^{-3/2} \text{ eV}. \quad (3)$$

For iron nuclei, equation \textbf{3} gives $E^b \lesssim 5 \times 10^{21} \Gamma \eta^{-1/2} (B/1 \text{G})^{-1/2} \text{ eV}$, which implies that in order to enable acceleration of cosmic rays to the highest observed energies, $\gtrsim 10^{20}$ eV, the strength of the magnetic field at the acceleration site should not exceed $B \lesssim$ few -few tens G (as long as $\Gamma \sim 1$, as is expected in RQ AGNs. See Ref. \textbf{57}).

The constraints of equations \textbf{1} \textbf{4} can be combined to constrain the radius of the acceleration site,

$$r_{\text{acc}} \equiv r \geq 9 \times 10^{16} \eta^2 \Gamma^{-2} Z A^{-4} E^b_{20}^{5/2} \text{ cm}. \quad (4)$$

For iron nuclei, equation \textbf{4} implies $r_{\text{acc}} \geq 6 \times 10^{19} \eta^2 \Gamma^{-2} \beta E^b_{20} \text{ cm}$, which is comparable to the Schwarzschild radius of typical black holes in AGNs (with characteristic mass $M_{BH} = 10^8 M_\odot$), $r_{\text{Sch,M}_\odot} = 3 \times 10^{13} \text{ cm}$.

Comparison of the synchrotron cooling time and the dynamical time shows that if the acceleration takes place at radii smaller than

$$r_{\text{acc}} \leq \frac{1.8 \times 10^{18}}{B^2} \Gamma^2 \beta (A/Z)^4 E^b_{20}^{-1} \text{ cm}, \quad (5)$$

then the synchrotron cooling time for the most energetic particles is shorter than the dynamical time, $t_{\text{cool, syn}}(E^b_{20} = 1) < t_{\text{dyn}}$. Assuming that the particle escape time is comparable to the dynamical time, when the condition in equation \textbf{5} is met, it implies that the spectrum of escaped particles is similar to that determined by the acceleration mechanism. We further discuss the implication of this condition in section \textbf{11} below. Combining equations \textbf{4} and \textbf{5} one obtains an upper limit on the value of the magnetic field at the acceleration site,

$$B \lesssim 4 \Gamma^2 A^4 Z^{-3} E^b_{20}^{-2} \eta^{-1} \text{ G}. \quad (6)$$

A lower limit on the value of the magnetic field can only be obtained once the acceleration radius $r$ is specified. This will be discussed in section \textbf{11} below.

In addition to synchrotron energy losses, energetic particles can in principle lose their energy by interacting with the ambient photon field and with other nuclei. Interaction with the photon field can result in energy losses
by Compton scattering (which results in negligible energy losses for UHECR’s, and therefore will not be considered here), photopair production (Bethe-Heitler process), photo-production of mesons (mainly pions), and photodisintegration of the nuclei. Therefore, the condition \( t_{\text{acc}} < t_{\text{cool}} \) as given in equation 2 should be modified.

A. Constraints from Disintegration of Energetic Nuclei

The threshold energy for photopair production, \( \epsilon_{\text{th}}^p = 2m_e(1 + m_e/m_p) \approx 1 \text{ MeV} \), is lower than the threshold energy for pion production, \( \epsilon_{\text{th}}^\pi = m_\pi(1 + m_\pi/2m_p) \approx 145 \text{ MeV} \). The relative contributions of these two processes to the energy loss of a relativistic particle is determined by the product of the cross section and inelasticity coefficient of the relevant process. For energetic protons, the product of cross section and inelasticity in pion production process is two orders of magnitude larger than the product in pair creation process \([66, 67]\). Exact calculation must take into account the target photon spectrum. It can be shown that for a photon number spectral index \( \alpha = 2 \) photomeson production is somewhat more dominant, however both time scales are comparable \([68]\). For heavy nuclei this ratio drops by a factor \( Z^2/\Lambda \) (see Refs. \([66, 69]\) and references therein), which for iron nuclei still implies that photomeson production is more important than photopair production as an energy loss channel.

However, the main energy loss channel of energetic particles is photodisintegration \([70]\). The threshold energy for this process is \( \epsilon_{\text{th}}^\gamma \approx 10 \text{ MeV} \), larger than the threshold energy for photopair production. Here, we thus focus on photomeson production and photodisintegration as the main energy loss channels of energetic particles interacting with the photon field.

An energetic nucleus having Lorentz factor \( \gamma_\Lambda \) (in the plasma comoving frame) propagating through an isotropic photon background with differential number density \( n(\varepsilon) d\varepsilon \) (at the energy range \( \varepsilon_\gamma \ldots \varepsilon_\gamma + d\varepsilon_\gamma \)) loses energy by photodisintegration and photomeson production at a rate (see, e.g., Refs. \([71, 72]\))

\[
\frac{1}{t_{\text{dis}, \gamma}} = \frac{c}{2\gamma_\Lambda^2} \int_{\epsilon_{\text{th}}}^{\epsilon_{\gamma}} \epsilon' \sigma_{A,\gamma}(\epsilon') \kappa(\epsilon') d\epsilon' \int_{\epsilon'/2\gamma_\Lambda}^{\infty} \frac{n(\epsilon'')}{\epsilon''^{2+\gamma_\Lambda}} d\epsilon'' \tag{7}
\]

Here, \( \sigma_{A,\gamma}(\epsilon') \) are the cross sections for photodisintegration and photomeson production, and \( \kappa(\epsilon') \) is the inelasticity coefficient in the photomeson production process. In calculating survival of one specie of heavy nuclei due to the photodisintegration process, equation 7 can still be used by inserting \( \kappa(\epsilon') \approx 1 \).

We estimate the differential photon number density at the acceleration radius \( r_{\text{acc}} \) in the following way: as a rough approximation, the spectrum of RQ AGNs increases as \( n L_\nu \propto \nu^2 \) at low energies, \( \nu < \nu_0^b \approx 10^{13} \text{ Hz} \), and is flat \( (n L_\nu \propto \nu^0) \) at higher frequencies, up to the X-rays (\( \nu_0^b \approx 10^{18} \text{ Hz} \); See Refs. \[73, 74, 75\]) \[96\]. The differential photon number density can therefore be written as a broken power law,

\[
n(\varepsilon) = n_0(\varepsilon/\varepsilon_0)^{-\alpha}, \tag{8}
\]

where \( \alpha = 0 \) for \( \varepsilon < \varepsilon_0 \) and \( \alpha = 2 \) for \( \varepsilon > \varepsilon_0 \). Here, \( \varepsilon_0 = h\nu_0^b/\Gamma = 6.6 \times 10^{-14} \Gamma^{-1} \text{ erg} \). The normalization constant, \( n_0 \) is found by normalizing to the total bolometric luminosity \( L_{\text{bol}} \equiv 10^{43}L_{\text{bol,43}} \text{ erg s}^{-1} \). Note that this bolometric luminosity does not correspond to the photon luminosity of the weak jets themselves, which are necessarily weaker. However, the survival of UHE nuclei depends on the entire photon field. Using the (comoving) energy density at the acceleration radius, \( u = L_{\text{bol}}/(4\pi r^2\Gamma^2\beta c) = \int n(\varepsilon) d\varepsilon \), one finds

\[
n_0 \approx \frac{4\pi c^2}{27\Gamma^2} \frac{\varepsilon_0}{\gamma_\Lambda^2} \left\{ \frac{1}{\beta c} \right\} \Gamma_{\text{bol}} \approx 6 \times 10^{10} L_{\text{bol,43}} \nu_0^{b,0.11} \text{ cm}^{-3} \text{ erg}^{-1}, \tag{9}
\]

where we approximated \( \log(\varepsilon_{\text{max}}/\varepsilon_0) \approx 10 \).

For the photon spectrum given by equation 8, the inner integral in equation 7 can be written as

\[
\int_{\epsilon'/2\gamma \Lambda}^{\infty} \frac{n(\epsilon'')}{\epsilon''^{2+\gamma_\Lambda}} d\epsilon'' = n_0(\varepsilon_0/\varepsilon_0) \left( \frac{\epsilon'}{(2\gamma_\Lambda \varepsilon_0)^{-\alpha}} \right)^{(1+\alpha)}, \tag{10}
\]

(Note that the above result is accurate for \( \epsilon'/2\gamma_\Lambda > \varepsilon_0 \), while for \( \epsilon'/2\gamma_\Lambda < \varepsilon_0 \) there is a second term that can be neglected for large enough values of \( \gamma_\Lambda \); see further discussion below.)

In order to estimate the outer integral in equation 7 we discriminate between photodisintegration and photomeson production processes. In the photodisintegration process, the main contribution to the outer integral in equation 7 is from photons in the energy bandwidth of the giant dipole resonance, whose energy is given by \( \epsilon_{\text{GDR}} = 42.65 \times A^{-0.21} \text{ MeV} \) for \( A > 4 \) \([76]\). Numerical fits to the experimental data gives the energy bandwidth of the resonance \( \Delta \epsilon_{\text{GDR}} = 8 \text{ MeV} \), and the maximum value of the cross section to be \( \sigma_{A,\gamma} = 1.454 \times 10^{-27} \text{ cm}^2 \). Approximating the outer integral by the contribution from the resonance, one finds

\[
\int_{c_{\text{min}}}^{\infty} \frac{d\epsilon' \sigma_{A,\gamma}(\epsilon') \mu_{\gamma}^0}{\epsilon_{\gamma}^{(1+\alpha)}} \equiv \frac{c_{\text{min}} \sigma_{A,\gamma} \Delta \epsilon_{\text{GDR}}}{(2\gamma_\Lambda \varepsilon_0)^{\alpha}} \tag{11}
\]

Note that this result is in good agreement with the more detailed numerical calculation at \( \gamma_\Lambda \approx \epsilon_{\text{GDR}}/2\varepsilon_0 \) (see Ref. \[14\] and figure \[17\] \[97\]).

The photodisintegration time can be compared to the dynamical time, \( \tau/\Gamma^3 \beta c \) \([23]\), to obtain a lower limit on the acceleration radius which allows survival of the accelerated nuclei. The result depends on the target photon spectrum. For the assumed spectral index \( \alpha \approx 2 \) above \( \nu_0^b \), the dissipation time obtains its minimum value at
For nuclei with this value of the Lorentz factor, one finds the survival condition,
\[ t_{\text{dyn}}^{-1} < 1 \rightarrow r \gtrsim 2.1 \times 10^{16} A^{1.21} (b_{\text{bol,43}}, b_{0.13})^{-2} \Gamma^{-4} \beta^{-2} \text{ cm}. \]

For iron nuclei, this gives \[ r \gtrsim 3 \times 10^{18} (b_{\text{bol,43}}, b_{0.13})^{-2} \Gamma^{-4} \beta^{-2} \text{ cm}. \] Therefore, for 10 parsec-scale weak jets (assuming here \( \beta \sim 0.3 \)) of RQ AGNs, we can expect that heavy nuclei can not only be accelerated but also survive without significant photodisintegration. We further note that for nuclei at higher energies, numerical calculations (e.g., Ref. [14]) indicate roughly comparable constraints on the acceleration radius. As at these high energies (\( \gtrsim \epsilon_{\text{GDR}}/2\epsilon_0 = 10^{19} \text{ eV} \)) the simple analytical treatment may be insufficient given the contribution from non-GDR resonance, we carried a numerical calculation. We present in figure 1 a comparison of the analytical approximation of the energy loss time to a more accurate numerical calculation, as well as the cooling time scales due to additional phenomena. While some deviation at the very high energies exist, clearly the analytical approximation used here is very accurate at lower energies, and does not deviate much from the exact numerical solution above \( 10^{19} \text{ eV} \) and below \( 10^{20} \text{ eV} \), for the parameters values chosen.

The second energy loss channel for energetic nucleons or nuclei is the photomeson production. The main contribution to the outer integral in equation (11) is from photons at energies \( \epsilon_{\text{peak}} \sim 0.3 \text{ GeV} \), where the cross section peaks at the \( \Delta \)-resonance, whose width is \( \Delta \epsilon \approx 0.2 \text{ GeV} \). For protons, the peak of the cross section at the \( \Delta \)-resonance is \( \sigma_{\text{peak}} \approx 5 \times 10^{-26} \text{ cm}^2 \), and the inelasticity is \( \kappa_{\text{peak}} \sim 0.2 \) (71). For heavier nuclei, the cross section is roughly proportional to \( A \) (71), while the inelasticity is proportional to \( A^{-1} \) (see Ref. [71]). Repeating the same calculation as for the photodisintegration process presented above, one finds that the limitation on the acceleration radius arises for particles at Lorentz factor \( \gamma_A \sim \epsilon_{\text{peak}}/2\epsilon_0 \). For these particles, we have \( t_{\gamma A}^{-1} \sim (2/(1+ \alpha))/n_0 \sigma_A \kappa_{\text{peak}} (\Delta \epsilon/ \epsilon_{\text{peak}})(\gamma_A/0.5 \epsilon_{\text{peak}})^{-1} \alpha^{-1} \) (98).

The requirement that \( t_{\gamma A}^{-1} < 1 \) leads to a much looser constraint than that for photodisintegration. Note that this condition does not depend on \( A \) for a given \( \gamma_A \). We thus conclude, that under conditions that enable survival of heavy nuclei, photomeson production is inefficient (see also Refs. [14, 69]).

The number density of particles at this distance, \( n \approx L/4 \pi r^2(2 \beta m_p c^3) \approx 2 \times 10^{-2} \rho_{\text{bol,43}} 2^{-2} \Gamma^{-2} \beta^{-1} \text{ cm}^{-3} \), implies that spallation is not important as an energy loss channel of the energetic particles. The rate for spallation process due to collision with other nuclei can be estimated as
\[ t_{\text{sp}}^{-1} = \sigma_{\text{sp}} n c, \]
where \( \sigma_{\text{sp}} = 5 \times 10^{-26} \text{ cm}^2 \) is the cross section for spallation of a nucleus (69). One can thus conclude that for acceleration at parsec scales, the spallation loss time is thus much longer than the dynamical time.

**III. ENERGETICS OF RQ AGNS AND EFFICIENCY OF UHECRS ACCELERATION**

In the previous section, we have concluded that the conditions inside the weak jets of RQ AGNs enable the acceleration and survival of UHE nuclei, if particle acceleration occurs at parsec scale radii. Such radii may be characteristic of internal or standing shocks in a jet, or shocks produced by interaction of the jet and the ambient medium. In this section, we compare the energy budget in RQ AGNs with the energy requirement for UHECRs acceleration. The observed flux of UHECRs on earth is \( J \approx 5 \times 10^{-18} \text{ m}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \) (see Ref. [2] and references therein). The production rate of UHECRs within the GZK horizon of \( R_{\text{GZK}} \approx 100 \text{ Mpc} \) is therefore \( \dot{N} = 4 \pi J/R_{\text{GZK}} \approx 1.7 \times 10^{36} \text{ Mpc}^{-3} \text{ yr}^{-1} \). For energies above \( 10^{19.5} \text{ eV} \), this implies energy production rate in energetic particles of \( \dot{E} \approx 5.5 \times 10^{43} \text{ erg Mpc}^{-3} \text{ yr}^{-1} \). This value is basically in agreement with more detailed estimates (see, e.g., Refs. [21, 48]).

Following Refs. [43, 57, 78] (and references therein), we consider a typical RQ AGN jet luminosity \( L_{\text{jet}} \sim 10^{42-43} \text{ erg s}^{-1} \). These values are based on measurements.
in the optical band \cite{78}: The typical radio emission is much weaker, \( \sim 10^{38-39} \text{ erg s}^{-1} \), due to both synchrotron self absorption and free-free absorption. Although the value for \( L_{\text{jet}} \) is highly uncertain, we find this luminosity to be equal to or somewhat fainter than the magnetic luminosity required for acceleration of UHE nuclei (say, with \( Z \gtrsim 10 \)), see eq. \[ \text{II} \]. The number density of RQ AGNs in the local universe is \( \phi_{42} \sim 10^{-3} \text{ Mpc}^{-3} \) which is about 10 times larger than RL AGNs (see, e.g., Ref. \cite{79}). Therefore, we expect that the energy input from RQ AGNs is \( \sim 3 \times 10^{46} \text{ erg Mpc}^{-3} \text{ yr}^{-1} \).

The number density of RQ AGNs drops with their luminosity \cite{80, 81}. For example, the number density of AGNs with soft X-ray luminosity, \( L_X \sim 10^{44} \text{ erg s}^{-1} \), is approximately \( \phi_{44} \sim 10^{-6} \text{ Mpc}^{-3} \). While at present the jet luminosity is not easily deduced from the X-ray luminosity measurements (e.g., absorption has a different effect on measurements at the soft and hard X band, see Ref. \cite{82}), we can still conclude that expected energy output from RQ AGNs jets significantly exceeds the energy requirement for UHECRs sources. Therefore, RQ AGNs jets are energetic enough to supply acceleration of particles to ultrahigh energies.

The estimated number density of RQ AGNs, \( \phi_{42} \sim 10^{-3} \text{ Mpc}^{-3} \), implies that for isotropic distribution of sources, the rate of energy production in UHE particles in a single source should be roughly \( \sim 3 \times 10^{40} \text{ erg s}^{-1} \) at \( 10^{19.5} \text{ eV} \), in order for these objects to be the sources of UHECRs (and assuming that no additional sources of UHECRs exist). This value is less than a percent of the expected typical jet luminosity of these objects.

This obtained value of \( \sim 10^{39.5} \text{ erg s}^{-1} \), while well within the energy limit of a typical RQ AGN, should be considered as a lower limit. First, the estimated number density of RQ AGNs does not take into account the possibility that in some of these objects the jet luminosities may be insufficient (see further discussion below). Second, the uncertainty that exists in the value of the efficiency parameter \( \eta \), as well as in the shock velocity \( \beta \), implies, via equation \[ \text{II} \] that the required luminosity for acceleration of iron nuclei to \( 10^{20} \text{ eV} \) may be somewhat higher than the value \( 10^{42} \text{ erg s}^{-1} \) considered here, in which case the number density of available sources may be lower. Therefore, even though our scenario remains essentially viable, one should keep in mind the high uncertainty that exists in the details of the physical processes. Nonetheless, the values obtained here can give a first approximation of the physical requirement from the source of UHECRs. The fact that they are well within the known energy budget of these objects leaves an ample margin to incorporate many of the uncertainties in the values of the physical parameters.

The above estimate does not take into account acceleration of particles to lower energies, which accompanies the production of UHECRs. Therefore, the required cosmic-ray input per source depends on the spectral index of the accelerated particles in the source. Although very steep source spectra can lead to an ‘energy crisis’ (see, e.g., Refs. \cite{21, 83}), RQ AGNs can be viable sources if sufficiently flat spectra are achieved. As both the jet luminosity and number density are uncertain at present, the results obtained here can perhaps serve as a guideline. Additional source of uncertainty lies in the energy losses caused by shock waves formed if a large number of charged particles escape from the source. While the energy of the escaping particles is reduced, still the highest energy particles, for which the escape time is the shortest, will be less affected. The exact details of this effect depend on the exact spectrum of the escaping particles, and are thus beyond the scope of this manuscript. However, we point here that if only the highest energy particles escape, this effect is not expected to play a significant role. Then, the overall cosmic-ray energy spectrum may be achieved as a superposition of contributions from many RQ AGNs with different cosmic-ray luminosities and maximum energies. Finally, our model requires a significant fraction of heavy element composition in the jet. As the jet material is ejected from the inner part of the disk, existence of metals in the disk inevitable leads to their existence in the jet, unless nuclei are disrupted in the jet base and the accretion flow in the disk.

The constraints found in \[ \text{II} \] above imply that UHECR production is possible if the acceleration radius is of the order of parsec, in order to avoid energy loss by photodisintegration. We note that the calculations carried in \[ \text{II} \] considered the bolometric luminosity, which is higher than the jet luminosity. As a result, the constraints derived above are more restrictive than the constraints that would have obtained by considering the (uncertain) luminosity in the jet only.

Using the results derived in \[ \text{II} \] we can obtain a lower limit on the value of the magnetic field at the acceleration site if we assume the acceleration radius \( r \), and characteristic velocity \( \beta \). If UHECRs are composed of iron nuclei, assuming \( \beta = 0.3 \), equation \[ \text{I} \] requires acceleration radius of \( r \sim 10 \text{ pc} \). This radius, the Hillas condition (equation \[ \text{I} \]) implies \( B \sim 10^{-4} \eta \beta^2 \frac{E^2}{L_{\text{bol,43}}} \). For carbon nuclei, the lightest nuclei that are consistent with the restrictions on the luminosity obtained in equation \[ \text{I} \] the minimum acceleration radius consistent with equation \[ \text{II} \] should be at \( 5 \times 10^{18} \text{ cm} \). The magnetic field at this radius should thus exceed \( B \geq 3 \times 10^{-3} \eta \beta^3 \frac{E^2}{L_{\text{bol,43}}} \). We note that for both scenarios, sub-Gauss values of the magnetic field is consistent with the requirement that the cooling time is shorter than the dynamical time (see equation \[ \text{III} \]), in which case the spectrum of the escaping particles is similar to that determined by the acceleration mechanism. We can thus constrain the required value of the magnetic field at the acceleration site to be in the range \( 10^{-3} - 1 \text{ G} \).

This \( \sim 1 \text{ parsec} \) acceleration radius derived here is consistent with the observed scale of weak jets observed in nearby Seyfert galaxies \cite{54, 55, 56, 57, 58}. At this radius, internal interactions with the jets, or interaction of the
jet with the interstellar medium material are expected to produce shock waves, that dissipate the jet kinetic energy. These shock waves are plausible acceleration sites for energetic particles. The lower limits on the values of the magnetic field derived above can be translated into lower limits on the equipartition value of the magnetic field at the acceleration site: using \( B^2/8\pi = \epsilon_B u \) (see discussion above equation 2), one finds that for iron nuclei, \( \epsilon_B \geq 4 \times 10^{-3} \), while for carbon nuclei, \( \epsilon_B \geq 10^{-1} \). These values are both below equipartition, and are consistent with values of the magnetic field expected to be produced in shock waves.

IV. OBSERVATIONAL IMPLICATIONS

In the previous sections, it has been shown that RQ AGNs can be viable sources of UHECRs, provided that the UHECRs composition is dominated by heavier nuclei (e.g., iron) rather than protons. This claim is based on two general considerations: (1) For a heavy nuclei composition, the required magnetic luminosity is weak enough to enable RQ AGNs to accelerate UHECRs; and (2) RQ AGNs are about ten times more numerous than RL AGNs, so that they can significantly contribute to the observed UHECR flux.

If the assumption presented here is correct, in addition to the contribution of RQ AGNs to the observed UHECR flux, we expect a contribution from RL AGNs. This is because the same acceleration mechanism can be expected to work in the latter. Although it is difficult to estimate the relative contribution of these two classes of objects, we expect that heavy nuclei (presumably, iron) from RQ AGNs can mask the correlation that may be found between the arrival direction of UHECRs and RL AGNs. In this respect, our scenario is consistent with the recent “disappointing model” \([52]\), which does not predict a significant correlation of the most energetic CR’s with AGNs.

Nonetheless, the conditions within the jets of RL AGNs may enable proton acceleration to ultra-high energies within these objects as well. Thus, if our scenario is correct, even if the correlation that can be expected between the arrival direction of UHECRs and positions of RL AGNs is reduced by the additional contribution from RQ AGNs, some correlation between the arrival directions of UHE protons and the positions of RL AGNs is still expected. The strength of this correlation depends on the efficiency of particle acceleration in RQ and RL AGNs, and on the composition of particle acceleration in RL AGNs: as our model allows, in principle, proton acceleration to UHE in RL AGNs, the observed UHECRs flux may be composed of two distinctive populations: heavy nuclei accelerated in RQ AGNs, and protons accelerated in RL AGNs. A testable consequence of this idea is that UHECRs whose arrival direction is not correlated with RL AGNs. Nonetheless, we stress that this is highly uncertain, due to two main reasons: first, RL AGNs may accelerate heavy nuclei to high energies as well; and second, the efficiencies of particle acceleration and escape from these sources are highly uncertain.

We can therefore conclude that our model does not rule out some correlation between the arrival direction of UHECRs and positions of RL AGNs. We note though, that the current observational status is highly uncertain. While some correlation was reported \([51, 57]\), its strength seems not to be fully determined yet (the recent results reported by Ref. \([37]\) indicate much weaker correlation than earlier reports \([31]\)). Therefore, future anisotropy search focusing on UHE proton events should be very important in confirming this idea.

In our scenario, the detection of secondary gamma rays and neutrinos from individual RQ AGNs whose luminosity is thought to be \( L \sim 10^{42–43} \) ergs\(^{-1}\) seems to be difficult. This is due to the fact that, as we saw in \([II A]\) in the weak jets in RQ AGNs photodisintegration is the dominant energy loss channel, hence copious production of energetic \( \pi^\pm \)'s is not expected. We leave the details of this calculation for future work.

On the other hand, the typical luminosity of jets of RL AGNs can be a few orders of magnitude brighter than RQ AGNs. Therefore, RL AGNs are more favorable targets for the purpose of detecting energetic secondaries. In particular, secondaries originating from very powerful AGNs like FR II galaxies (see, e.g., Ref. \([84]\) and references therein) or very nearby RL AGNs such as Cen A (e.g., Ref. \([23]\) and references therein) might be detected by Fermi or IACTs. Clusters of galaxies hosting powerful AGNs might also be viable candidate for secondary detection surveys (see, e.g., Refs. \([83, 85]\) and references therein).

V. SUMMARY AND DISCUSSIONS

In this paper we have considered radio quiet AGN’s as possible sources of UHECRs. So far, jets in these objects were not considered as sources of the highest energy cosmic rays around \( 10^{20} \) eV, since they are not luminous enough to support acceleration of protons to these energies. However, the recent results of the AUGER collaboration shows indications that at the highest energies the composition of cosmic rays may be dominated by heavy nuclei \([38, 39]\). The assumption that UHECRs are heavy nuclei cases the constraint on the source luminosity (see eq. 2). Thus it allows, in principle, radio-quiet AGNs to be the sources of UHECRs.

We have calculated in \([II A]\) the constraints on the acceleration site which enable the acceleration of heavy nuclei to high energies. The most restrictive constraint arises from the photo-disintegration process. We showed in \([II A]\) that for typical nearby AGN’s with bolometric luminosity \( L_{bol} \approx 10^{43} \) erg s\(^{-1}\), energetic nucleus can sur-
vive photo-disintegration if the acceleration takes place on a parsec scale (see eqs. [12]). Interestingly, this is the same scale at which weak jets were seen in these objects \[54, 55, 56, 58\]. This fact further supports our idea, since interaction of jets with the surrounding material inevitably leads to creation of shock waves, which are the most plausible acceleration site of particles to high energies.

The question of the acceleration of particles to ultra-high energies may depend on the shock velocity. In the Bohm approximation $\eta \propto \beta^{-2}$, some assumptions about the shock velocities in the jets are needed. For example, a shock velocity of $\beta \sim 0.1$ requires magnetic luminosity of $\epsilon_B L \gtrsim 10^{43.5}$ erg s$^{-1}$ in order to enable acceleration of iron nuclei to the maximum observed energy (see §II, equation [2]), which may be too low. However, somewhat higher values, $\beta \sim 0.3$ may be sufficient. Since currently there is a high uncertainty in the determination of the shock velocities in RQ AGNs, it is not possible to put further constraints. We have also showed in §III that the minimum inferred values of the magnetic field required to confine the accelerated particles at parsec scale jets is sub-Gauss. This value was shown to be smaller than the equipartition value, which is also consistent with models of particle acceleration in shock waves that may be produced on this scale.

We further showed in §III that the required luminosity in UHECRs could be less than a percent of the total bolometric luminosity of more abundant nearby AGNs. If the spectrum of produced cosmic rays is a power law with index close to $d \log N / d \log E \approx -2$, as suggested by models of particle acceleration in non-relativistic shock waves \[62\], then the total luminosity in energetic particles (at all energies) could still be much smaller than the total bolometric luminosity. Even for higher power law index, $p \approx 2.3$, we find that the total energy requirement for acceleration of cosmic rays above GeV is $\approx (10^{40})^{0.3} = 10^3$ times higher than the energetic requirement from UHECRs alone, which is still (marginally) consistent with the total energy budget in RQ AGNs (see discussion in §III). Assuming that the main radiative source in these objects is the accretion disk, this implies a high efficiency, of the order of tens of percents, in the conversion of accretion energy to acceleration of particles. Since large uncertainties exist in the efficiencies of both the energy conversion in the jet and the acceleration of particles, we can only conclude that our model is consistent with the energetic requirements, provided a high efficiency is achieved in both these processes.

Spectral synthesis as well as chemical enrichment models predict that AGN’s are metal rich. The metallicities in the broad emission line region is typically $\sim 1$ to $\gtrsim 10$ times the solar metallicity \[54, 57\], and grows with the luminosity \[54]. Thus, abundant existence of heavy nuclei in AGN’s disks is expected. As the plasma jet is composed of material from the disk, abundant population of heavy nuclei in the jet is expected. Thus, AGN’s may be one of the natural sources of high energy heavy nuclei.

If indeed the acceleration takes place on a parsec scale, then the main energy loss channel is photodisintegration, rather than photomeson production. As discussed in §IV, this fact implies that we do not expect abundant production of energetic neutrinos, that may result from the decay of energetic pions. On the contrary: the results in equations [12] indicate that the relative contribution of photomeson production may be no more than few percents of photodisintegration as an energy loss channel of UHECRs. Thus, we do not expect a copious production of neutrinos under the conditions which enable the acceleration and survival of heavy nuclei, as considered in this paper.

After they escape from their sources, UHECR nuclei can be subject to photodisintegration in the intergalactic space. UHE nuclei with energy $\gtrsim 10^{19}$ eV mainly interact with the cosmic infrared background (CIB) photons \[50, 51, 89\]. UHE iron nuclei with energy $\gtrsim 10^{20}$ eV have a mean free path of $\gtrsim 500$ Mpc, while the mean free path of oxygen nuclei at a similar energy is only $\sim 30$ Mpc \[60\]. Both increase rapidly as the energy of the nuclei decreases. Thus, the GZK horizon of UHE nuclei is comparable to, or even somewhat larger than the GZK horizon of energetic protons (see also Ref. [1]).

To summarize, we have pointed out here that the assumption that UHECRs are composed of heavy nuclei, as is suggested by recent AUGER results \[38, 39\], enables radio quiet AGN’s to be their sources. This picture is consistent with what is currently known about the existence of parsec scale jets seen in some of these objects, as well as the energy constraints. Moreover, this picture is supported by the recent analysis carried by Ref. [91], in which a strong correlation between the arrival directions of UHECRs and the sky coordinates of AGN’s detected by the Swift-BAT within the GZK cutoff (which are largely radio-quiet) was found. As more data are collected by the AUGER collaboration, a clearer picture of the composition of UHECR’s will become available, allowing further constraints on possible sources.

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Note though that [53] showed recently that radio emission from jets may be suppressed if the emitting particles cool rapidly enough close to the base of the jet, so that synchrotron emission from outer parts of the jet is self absorbed. Although this work was done in the context of jets in microquasars, a similar idea can be easily implemented in the case of RQ AGNs as well. However, a general discussion on the role of synchrotron-self absorption on the distinguishing between RL and RQ AGN’s is outside the scope of this paper, and is left for future work.

Note that we adopt here the common assumption that the particle velocity in the shock frame is comparable to the shock velocity.

Although we may not expect relativistic motion in the weak jets of radio quite AGNs, so that \( \Gamma \gtrsim 1 \), we introduce this quantity here for completeness.

Additional possible solutions to the problem of high luminosity presented in equation \([2]\) is obtained either (1) by assuming that the energy density in the magnetic field at the acceleration site is larger by 1-2 orders of magnitude than the energy density in the radiating electrons. In such a scenario, acceleration of protons to the maximum observed energy of \( 10^{20} \text{ eV} \) is also acceptable. (2) A second possibility is large luminosity flaring activity, as might have been seen by comparing \textit{FERMI} to earlier \textit{EGRET} observations [64].

The actual spectrum is of course much more complicated than this simple estimate. The radio emission may originate from weak jets, infrared emission may come from AGN tori while UV and x-ray emission are typically radiated by the accretion disks and coronae. The estimation given here is therefore mainly valid for checking the possibility of survival of UHE nuclei, but may not be valid for more accurate predictions of hadronic emission.

The contribution from the non-GDR-resonance (including quasi-deuteron and fragmentation processes) region becomes important at \( \gamma_A \lesssim \epsilon_{\text{GDR}}/2\epsilon_0 \) for \( \alpha \lesssim 1 \).

More accurately, the photodisintegration time should be compared to the escape time scale, which is highly uncertain. If the escape time scale is too long, UHECRs would lose their energy due to, e.g., the adiabatic cooling (whose time scale is typically \( t_{\text{ad}} \sim 3t_{\text{dyn}} \)). Here we assume that \( t_{\text{esc}} \) is longer than but comparable to \( t_{\text{dyn}} \), hence UHECRs can escape from the source without significant adiabatic loss.

In a similar way to photodisintegration process, at high energies \( \gamma_A \gtrsim \epsilon_{\text{peak}}/2\epsilon_0 \), the energy loss time differs from the simple estimate given here if \( \alpha \lesssim 1 \) due to multi pion production. However, the \( \Delta \)-resonance approximation used here is still valid at \( \gamma_A \lesssim \epsilon_{\text{peak}}/2\epsilon_0 \).

Note that similar distances are obtained for proton as well heavy nuclei composition of UHECRs. See, e.g., Ref. [1].