Gamma rays and neutrinos from a powerful cosmic accelerator

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Abstract

Possibly, the powerful radio quasar 3C273 will reveal its nature as an efficient proton accelerator up to energies of order $10^{11}$ GeV in the near future. It is shown in this paper that the shock accelerated protons expected to be present in the quasar’s plasma jet induce an unsaturated synchrotron cascade with electromagnetic radiation emerging in the X–ray and gamma ray range. While (including the synchrotron emission from the accelerated primary electrons) the broadband nonthermal emission from 3C273 can be explained over the observed 18 orders of magnitude, a flattening of the spectrum at the highest observed energies (a few GeV) is predicted that could be falsified by the Energetic Gamma Ray Experiment Telescope (EGRET) on board the Compton Gamma Ray Observatory (GRO). Above $\approx 100$ GeV the cascade spectrum dramatically steepens again due to the absorption of the gamma ray photons.
by the host galaxy’s strong infrared photon field from extended dust clouds – in accordance with the non-detection of 3C273 by Cherénkov telescopes. However, neutrinos from the hadronic interactions initiating the cascade are not damped and reach terrestrial experiments without any modification of their injected flux. In contrast to the neutrino flux from pp–interactions, which are energetically unimportant in jets, pγ–interactions generate a flat neutrino flux. Therefore it is emphasised that one must not simply normalize the expected neutrino flux by the observed gamma ray flux. Hence it is shown that the expected neutrino flux in the energy range relevant for underwater or underice detectors is much lower than assumed by many authors. On the other hand, with an increasing number of cosmic gamma ray sources at known positions, their neutrino detection should be feasible when it is realized that angular resolution is the crucial design property for neutrino detectors.

95.30, 95.70.R, 98.70.S, 96.40.P
I. INTRODUCTION

Recent gamma ray observations by GRO may be the clue to the long-standing problem of the origin of UHE-cosmic ray protons. GRO has discovered that many flat-spectrum radio quasars are strong gamma ray sources with energy spectral index \( s = \log [dN_\gamma /dE]/d \log [E] \) above 100 MeV ranging between roughly \(-1.5\) and \(-2.5\). As pointed out below the emission indicates the presence of very energetic baryons that interact with the dense photon atmosphere in the plasma outflow emerging from the active galactic nucleus (AGN). A prominent example is the nearby quasar 3C 273 with index \(-2.39 \pm 0.13\) which is investigated in this paper because the steep gamma ray spectrum observed may help to discriminate between the proton powered and other – purely leptonic – models.

A. The nonthermal “mini-blazar”

3C 273 is a well-known superluminal radio quasar at redshift \( z = 0.158 \) (e.g. Ref. 2). VLBI experiments showed that radio knots propagate along a strongly collimated jet emerging from a stationary core with apparent opening angle \( \Phi_{ob} = 4^\circ \). The motion of the knots with respect to the core is superluminal with apparent speed \( \beta_{ob} = 7h^{-1} \) over a projected distance of at least \( r_{ob} = 45h^{-1} \) pc at 5 GHz where \( h = H_0/100 \) Mpc km \(^{-1}\) s \(^{-1}\).

The total flux of the jet comprises the flux of several knots with increasing self-absorption frequency towards the base of the jet and is flat \( (\alpha \simeq 0) \) up to a spectral break in the mm range above which the spectrum steepens \( (\alpha \simeq 1) \). Detailed studies of the variability and polarization of 3C 273 have disentangled the nonthermal emission component associated to the compact jet (the low-entropy “miniblazar”) from the thermally emitting (high-entropy) sites, such as the UV and soft X-ray excess emitting “cold” gas or the infrared emitting dust.

The miniblazar can be attributed to a relativistic plasma outflow from the AGN with embedded magnetic fields and shocks. The shocks accelerate particles and the accelerated
electrons lose their energy rapidly due to observable radio-optical synchrotron emission in the magnetic field. Interferometric radio maps rather suggest a helix pattern for the trajectories of the radio knots and, indeed, the situation can be pictured as a topological equivalent to the curved shocks in the solar wind following the Parker spiral of the magnetic field – with the equatorial plane of the sun being bent onto a cone, so that the spiral becomes a helix, and the shock waves being excited by mass ejections from the accreting gas at the base of the jet. Remarkably, the accelerated protons and nuclei were long thought to be negligible as radiative agents, since they have much lower energy losses at the same energy as electrons. However, acceleration thus operates longer for the heavy particles until their energy losses also become large. The energy lost by the energetic baryons then emerges as cascade radiation in the X-gamma ray band.

Since the shock speed does not reflect the fluid speed (in which the scattering centers move with the Alfvén speed), the speeds of knots may be different, i.e. a strong reconfinement shock after an initial passage of a thick cloud may be a standing or a weakly–relativistic shock wave. As a consequence the emission from this shock, particularly the infrared/X–ray emission, could be less beamed than the radio emission, since radio emission is self–absorbed within the inner parsec of a jet and thus comes from further out. Beaming of the emitted radiation pattern results from the anisotropy arising from the Lorentz transformation between the comoving fluid frame (with isotropic particle and photon distributions) and the observer’s frame (“Doppler–boosting”). However, for the sake of simplicity it is assumed in this paper that the Doppler factor is the same along the jet.

B. UHE–protons

The crucial question is, of course, whether protons can really reach the extremely high energies needed to generate a luminosity of secondaries competitive with inverse–Compton or synchrotron–self–Compton emission from the accelerated electrons and to be able to provide enough energy in the center–of–momentum frame of the colliding protons and photons to
create secondary particles on the mass shell.

The maximum possible energy is readily obtained from

$$t_{p,\text{cool}}(E_{p,\text{max}}) = t_{\text{acc}}(E_{p,\text{max}})$$  \hspace{1cm} (1.1)

provided that the limit from drift across the shock with radius $r_\perp$

$$E_{p,\text{max}} \approx e (v \times B) r / c \simeq 7.8 \cdot 10^{20} \beta_s \left( \frac{B}{G} \right) \left( \frac{r_\perp}{\text{pc}} \right) \text{eV}$$  \hspace{1cm} (1.2)

is not exceeded. Here $t_i$ denotes the time scale of process i, $B$ the magnetic field and $\beta_s$ the speed of the shock. Due to energy losses the proton distribution steepens from $dN/dE \propto E^{-2}$ (or even flatter, cf. Ref. [8]) to $dN/dE \propto E^{-3}$ when

$$t_{p,\text{cool}}(E_{p,b}) = t_{e,\text{cool}}(E_{e,b}) = t_{\text{exp}}$$  \hspace{1cm} (1.3)

where $t_{\text{exp}}$ denotes the dynamical time scale of the expanding jet. When $E_{p,b} \geq E_{p,\text{max}}$, no steepening occurs and the proton distribution turns over with a Bessel–function type behaviour at $E_{p,\text{max}}$. Since

$$\frac{L_p}{L_e} \simeq \frac{u_p}{u_e} \frac{t_{e,\text{cool}}}{t_{p,\text{cool}}} \simeq \eta \frac{\text{Min}[E_{p,b}, E_{p,\text{max}}]}{E_{p,b}}$$  \hspace{1cm} (1.4)

a small $E_{p,\text{max}}$ costs a great value of $\eta = u_p/u_e$ to obtain comparable luminosities from protons and electrons, respectively. Note that $\eta$ in Eq. (1.4) refers to the energy density ratio at extremely relativistic energies, so that the Galactic value $\eta(\geq \text{GeV}) \approx 100$ indeed lets one expect the possibility that $L_\gamma \gg L_{\text{ir}}$. With $B \propto r^{-1}$ along the jet of length $r$ and $r$ small enough, so that $E_{p,b} \leq E_{p,\text{max}}$, it follows that $L_p/L_e \propto \eta$, whereas further out $L_p/L_e \ll \eta$. Therefore, one expects to observe gamma rays mainly from the compact regions of radio jets close to the core. Additionally, the emission from the nuclear jet (the blazar) is concentrated in a narrow lighthouse beam due to relativistic bulk motion of the plasma amplifying the flux at small angles to the line of sight. Thus it can be explained, why flat–spectrum radio sources (indicating jet sources at small angles to the l.o.s.) are the gamma ray sources detected in flux–limited experiments.
Comparing 3C 279 and 3C 273 one can see that in spite of a similar inferred radiation compactness, the maximum proton energy in the jet of 3C 273 is smaller because of the presence of the strong external radiation field seen as the big blue bump. For the maximum dimensionless energy it follows from Eq. (1.3) that

\[ \gamma_{p,b} = \left( \frac{m_p}{m_e} \right)^3 \frac{1 + a_s}{1 + \langle \sigma_{p\gamma}/\sigma_{p,\text{syn}} \rangle a_s} \gamma_{e,b} \]  

with the target photon/magnetic energy density ratio \( a_s = u_\gamma/u_B \) and the factor \( \langle \sigma_{p\gamma}/\sigma_{p,\text{syn}} \rangle \approx 240 \) accounting for the relative importance of photoproduction and synchrotron energy losses, respectively. For the intrinsic photons \( a_{s,\text{in}} \approx \beta_j \gamma_j \Phi/(1 + \eta) \ll 1 \), whereas the external photons with luminosity \( L_{\text{uv}} \approx 10^{47} \) erg/s enhance the local photon density such that \( \alpha_{s,\text{ex}} = 2L_{\text{uv}}/(r_b + r_c)^2 B_b^2 c \approx 0.1 \) assuming \( r_b = r_c \) for the length \( r_b \) of the radiative jet and the distance \( r_c \) of the base of the jet to the location of the thermal photons (i.e. collimation out to the parsec scale). The value thus obtained for the proton break energy \( \gamma_{p,b} = 3 \cdot 10^{10} \) (Tab. 4) is in modest agreement with the value for \( \gamma_{p,\text{max}} \) obtained from Eq. (1.1). Thus, head-on collisions with a (infrared) photon of energy \( \varepsilon_t = 145 \text{ MeV}/2\gamma_{p,b} \approx 2 \cdot 10^{-3} \) eV can excite the \( \Delta \) resonance. Note that the observed energies are cosmologically redshifted and Doppler blueshifted yielding \( E_{p,\text{max}}(\text{obs}) = D_j \gamma_{p,\text{max}}/(1 + z) \approx 2 \cdot 10^{11} \) GeV, if there were no energy loss during propagation through the cosmic microwave background.

Since the pair absorption cross section peaks for head-on collisions, but the external photons appear highly anisotropic in the comoving fluid frame, the pair creation opacity (\( \propto \) compactness) in the very forward direction remains unaffected by the additional target photons. In contrast, the Thomson scattered luminosity is found to be \( L_T \approx a_{s,\text{ex}} L_s \), where \( L_s \) denotes the apparent soft synchrotron luminosity of the jet, since the Thomson scattering is isotropic in the comoving frame. In models where the particle acceleration takes place much closer to the AGN, i.e. \( r_c \ll 1 \) pc, the thermal seed photon energy density could exceed the magnetic field energy density in the jet, so that, indeed, the Compton–upscttered thermal photon flux can dominate over the emission from the jet itself (cf. Ref. [10] or for another
version of the same model see Ref. [11]). Very close to the AGN the extremely strong Compton cooling is faster than diffusive acceleration processes which are rather slow for reasonable diffusion coefficients. Only explosive acceleration seems possible in this case. On the other hand the well–ordered structure of the jet on much larger scales must not be destroyed by such processes, nor must there be too efficient dissipation, for otherwise the luminosity of the large scale jet could not be as large as observed.

In the following Sections it is shown how the cascade responds to a varying proton maximum energy and how the infrared photons from dust clouds modify the emerging spectrum. Finally, the neutrino flux from 3C273 is calculated.

II. THE UNSATURATED SYNCHROTRON CASCADE AND THE FEW MEV BUMP

It is instructive to calculate step by step the development of the unsaturated synchrotron cascade induced by the UHE protons. For each step conservation of the total emitted power \( L \propto x^2 dN_\gamma / dx \) determines the normalization of the stationary photon spectrum \( dN_\gamma / dx \).

The result of a numerical calculation is shown in Fig. 1 for the physical conditions listed in Tab. I and the detailed integral equation solved is discussed in [12].

We start with protons of Lorentz factor \( \gamma_{p,b} = 3 \cdot 10^{10} \) producing pions of Lorentz factor \( \gamma_\pi = (m_p / m_\pi) \gamma_{p,b} / 5 \). The neutral pions decay yielding two gamma rays of dimensionless energy

\[
x_0 = \frac{E_{\gamma}}{m_ec^2} \approx \frac{m_p}{m_e} \frac{1}{10} \gamma_{p,b} \approx 180 \gamma_{p,b}
\]

in the comoving frame (multiplication with the redshifted Doppler factor \( D_j / (1 + z) \) yields the observed frequencies or energies). The local emissivity of gamma rays is \( Q_\gamma \propto x^{-1} \) and the stationary photon spectrum steepens to \( dN_\gamma / dx \propto x^{-2} \) because of the \( \gamma\gamma \) pair production losses. Here we take \( dN_p / d\gamma_p \propto \gamma_p^{-2} \) (shock acceleration, negligible losses below \( \gamma_{p,b} \)) and \( dN_{\gamma,\text{target}} / dx \propto x^{-2} \) (as observed between \( 10^{11} \) Hz and \( 10^{15} \) Hz corresponding
to proton energies at threshold between $\gamma_1 = m_\pi/(2m_e x_1) = 10^{12}$ and $\gamma_2 = 10^8$, resp.).

The reason for the flat injected flux of secondary gamma rays is the following: Compared to cooling on monoenergetic target photons with energy $x_m$, which yield $Q_\gamma \propto x^{-2}$ above $\gamma_m = m_\pi/(2m_e x_m)$, cooling in the inverse power law photon distribution with $N_{\gamma,\text{target}}(x) \propto x^{-1}$ yields additional target photons ($x < x_m$) for $\gamma_p > \gamma_m$ and hence a flatter emissivity $Q_\gamma \propto x^{-2+1}$ [13]. Now, we forget about charged pions, direct pair production and proton synchrotron emission, because the pairs from charged pion decay simply add more flux to the neutral pion decay cascade, direct pairs radiate mostly around TeV where – as will be shown – the emission is absorbed by dust infrared emission and the proton synchrotron emission contributes only for $a_s < 0.004$.

The first synchrotron photon generation is now radiated by the pairs produced from the original neutral pion decay quanta via

$$\gamma + \gamma_t \rightarrow e^+ + e^-$$

where $\gamma_t$ denotes a soft photon at threshold energy $x x_t = 1$ (head–on collision). The maximum electron Lorentz factor is given by $\gamma_1 = x_o/2$ and the characteristic synchrotron frequency is then

$$x_1 = 3 \cdot 10^{-14} B_\perp \gamma_1^2 = 2.2 \cdot 10^{11} B_\perp$$

The stationary electron distribution being subjected to synchrotron losses is $dN_e d\gamma_e \propto \gamma^{-2}$ (for $\gamma < \gamma_1$) and therefore the synchrotron emissivity is $Q_\gamma \propto x^{-3/2}$ (for $x < x_1$). As long as we are in the optical thick range, the stationary photon spectrum is steeper by one power, i.e. $dN_\gamma/dx \propto x^{-5/2}$. The pairs thus produced have Lorentz factors $\gamma_2 = x_1/2$ and radiate the second synchrotron generation

$$x_2 = 3 \cdot 10^{-14} B_\perp \gamma_2^2 = 3.6 \cdot 10^8 B_\perp^3$$

with stationary photon spectrum $dN_\gamma/dx \propto x^{-11/4}$. However, part of this generation is already at energies below
\[ x^* \simeq \frac{10}{l} = 5 \cdot 10^6 \]  

(2.4)

where the jet becomes optically thin (compactness in the comoving frame \( l = 2 \cdot 10^{-6} \) for 3C 273. cf. Tab.1), and therefore has a spectrum \( dN_\gamma/dx \propto x^{-7/4} \). The energy \( x^{**} \) where the pairs produced at \( x^* \) radiate is given by

\[ x^{**} = 3 \cdot 10^{-14} B_\perp \left[ \frac{x^*}{2} \right]^2 = 0.2B_\perp \]  

(2.5)

The third generation has the characteristic energy

\[ x_3 = 3 \cdot 10^{-14} B_\perp \gamma_3^2 = 10^3 B_\perp^7 \]  

(2.6)

Now, here comes the most important point: There are two special energies involved. The energy \( x^* \) defined by \( \tau_{\gamma\gamma}(x^*) = 1 \) and the energy \( x_2 \), which may be either \( <, = \) or \( > \) than \( x^* \). In the first case the second cascade generation is optically thin with respect to pair creation (no further reprocessing) and has a spectrum with index \(-7/4\). In the second and in the third case, there will be a third cascade generation. The part of the spectrum from the optically thick range \( x_2 > x > x^* \) (which is monoenergetic for \( x_2 = x^* \)) has the form \( dN_\gamma/dx \propto x^{-15/8} \) over an energy band \( \Delta \log \left[ x \right] = 2 \log \left[ x_2/x^* \right] \) (a lá 3C 279) from \( x^{**} \) to \( x_3 \) and the part from the optically thin range has the spectrum \( dN_\gamma/dx \propto x^{-11/8} \) emerging below \( x^{**} \). Thus, \( \alpha_\gamma \simeq 0.9 \) and \( \alpha_X \simeq 0.4 \) ignoring that the superposition of individual cascade generations somewhat smears out these values.

The case \( x^* \approx x_2 \), relevant for 3C 273 where \( x_2 \) follows the reduction of \( \gamma_{p,b} \), leads to a gap between the maximum of the third generation at \( x^{**} \approx x_3 \) and the maximum of the second (partially absorbed) generation at \( x^* \). It is not possible to predict the accurate shape of this gap, because it depends sensitively on the upper turnover of the injection modulo all decay kinematics and successive synchrotron smearing (\( x \propto \gamma^2 \)). The numerical result (assuming \( B_\perp \approx B_b \) so that \( D_j x_3(1 + z) \approx 2 \), that is 1 MeV) with injection of an exponential turnover is shown in Fig.1 for 3C 273 yielding an effective spectral index close to the observed \(-2.39\). However, at a few GeV the spectrum of the decadic power rises again due to the onset of the second cascade generation with photon index \(-7/4\) up to 10 TeV \( \approx D_j x^*/(1 + z) \).
The energy \( x_3 \propto B^7 \perp \) is extremely dependent upon the magnetic field strength in contrast to \( x^{**} \propto B \perp \). Therefore one may expect a wide range of spectra in the MeV–GeV range. There is no way, however, to understand a few MeV bump without an accompanied intrinsic bump at 1–10 TeV, since the cascade generation preceding the one that makes the MeV bump must show up at \( D_j x^*/(1 + z) \).

### III. INFRARED ABSORPTION OF THE FEW TEV BUMP

This brings us to the final argument, viz. the absorption of gamma rays outside of the jet by photons from the extended regions of 3C 273.

The optical depth for scattering on a monoenergetic photon target is approximately given by

\[
\tau_{\gamma\gamma}(x) = r_{\text{ext}} N_\gamma(x_o) \sigma(x)
\]  

(3.1)

where \( x = h\nu/m_e c^2 \) denotes the dimensionless gamma ray photon energy, \( \sigma(x) \approx 3 \sigma_T x_{\text{th}}/x \) for \( x \geq x_{\text{th}} = 2/[x_o(1 - \cos \theta)] \) is the pair creation cross section above threshold energy \( x_{\text{th}} \), \( \sigma_T \) the Thomson cross section and \( N_\gamma(x_o) = L_o/(4\pi r_{\text{ext}}^2 m_e c^3 x_o) \) the external target photon density of the source with monochromatic luminosity \( L_o \) at target energy \( x_o \). For \( L_o = 10^{46} L_{46} \) erg/s and \( x \geq x_{\text{th}} \) this yields

\[
\tau_{\gamma\gamma}(x) = 4 \cdot 10^{15} L_{46} (r_{\text{ext}}/\text{cm})^{-1} x_{\text{th}}^{-1} x_o^{-1}
\]  

(3.2)

so that at threshold \( \tau_{\gamma\gamma} = 4 \cdot 10^{15} L_{46} (r_{\text{ext}}/\text{cm})^{-1} x_o^{-1} \). Assuming that the infrared emission comes from dust at \( r_{\text{ext}} = 300 \) pc we obtain that \( \tau_{\gamma\gamma} \geq 1 \) for \( x \geq 5 \cdot 10^5 \) corresponding to \( E_\gamma \geq 100 \) GeV. Since the produced pairs would isotropize rapidly, the reemitted radiation in the m.f. of the host galaxy does not remember the beam, so that the apparent luminosity is reduced to the comoving frame luminosity which is down by a large kinematic factor.

One can further ask, wether the flattening of the predicted cascade spectrum above 10 GeV (cf. Fig. [1]) could be destroyed by photon–photon absorption as well as the 10 TeV
bump. The absorbing thermal optical/X-ray photons presumably come from behind the jet (from the apex) and therefore the angle $\theta$ should have the value obtained for the angle between jet axis and observer, viz. $\theta = 7^\circ$. With $x$ denoting the observed gamma ray energy, the resonant target photon energy is $x_0 \simeq 1.4 \cdot 10^{-2}$ (7 keV). The optical depth at threshold for $L_{46}(7\text{keV}) = 0.1$ is given by $\tau_{\gamma\gamma}(r_c + r_b = 8\,\text{pc}, \theta = 7^\circ) \simeq 6 \cdot 10^{-4}$; clearly insufficient to steepen the spectrum.

IV. CONCLUSIONS

Radio jets emerge out of the vicinity of an accreting compact object, presumably a black hole, where a rotating magnetosphere tied to the accreting gas feeds a considerable fraction of the infalling matter into the outflow, thereby transporting angular momentum outwards. The magnetic fields allow for collimated flow solutions [14] passing through bulk Lorentz factors of 5-20 beyond the lightcylinder at roughly $10r_G \simeq 1.5 \cdot 10^{15}m_9 \text{ cm}$, where $m_9$ denotes the mass of the black hole in units of $10^9 m_\odot$. If the initial acceleration of plasma would be much more efficient, i.e. $\gamma_i \approx 10^3$, the ambient thermal UV and soft X-ray photons from the accreted gas exert a ‘Compton drag’ on the particles in the outflow slowing it down to terminal Lorentz factors $\gamma_\infty \approx 10$, thus dissipating most of its energy into Compton-scattered radiation [10]. However, this produces maximum photon energies well below 100 MeV – in apparent contradiction with the observed fluxes above GeV from many extragalactic radio sources. Moreover, if radio jets consist of ordinary hydrogen plasma, radio-loud objects like 3C 273 have a very powerful jet even many kiloparsecs away from the AGN and therefore it seems unlikely that they have dissipated much of their power in the central parsec already.

After some initial collimation by toroidal fields the jet expands freely and super-Alfvénically into the surrounding medium. However, free expansion tends to bring the internal pressure of the jet within a few scale heights above the Alfvén point into equilibrium with the external pressure. The flow tries to realise an equilibrium state by reconfinement
shocks, which let the internal pressure vary in a zig–zag fashion about the external pressure. Thus, the external gas determines the structure of the jet and the amount of radiative dissipation, which establishes itself through shock acceleration at the reconfinement edges. Another important source of shocks could be mass ejection at the base of the jet – much like in the solar wind. Thus, when the ejection sites are tied to a rotating accretion disk, a helical shock pattern results, which seems to be required by VLBI observations of many radio jets and HST observations of the M87 jet.

Synchrotron emission of accelerated electrons emerges mostly in the radio to optical regime, with the submm/infrared/optical emission coming from the innermost nuclear jet \((r < r_b \approx 4 \text{ pc})\) and the radio emission coming from further out. Note, however, that the propagation of the jet through the central parsec (out to the BLR) is assumed to be essentially non–dissipative, \(r_{c} = O[1\text{pc}]\). During this collimated stage the initial Poynting flux is converted into bulk kinetic flux, which is then the reservoir for particle acceleration at shocks farther out. Higher frequencies than optical are difficult to achieve from synchrotron emission by the primary accelerated electrons, because the maximum energies are strongly constrained by energy losses. However, the protons accelerated in the jet reach energies so high \((\gamma_{p,\text{max}} \approx 10^{10})\) that saturation effects due to the finite size of the shocks become important, simply because their cooling time scale, which is dominated by secondary particle production for a wide parameter range, is very long compared to the cooling time scale of electrons. Thermal photons from outside of the jet can additionally damp the protons in the jet. The power induced by the cooling ultra–high–energy protons at comoving frame energies up to \(10^{10} \text{ GeV}\) is rapidly reprocessed by an electromagnetic shower and emerges in the X– to gamma ray regime. The infrared to optical emission from the throat of the jet serves as the scattering agent for the cascade determining the photon energy \(\epsilon_{\gamma} \approx 10 \text{ TeV}\) where the jet becomes optically thick with respect to pair creation in the observer’s frame.

When the initial jet collimation terminates very far away from the central engine or when the thermal emission from the central engine is inherently weak, then the emerging cascade spectrum is much like as it is observed by GINGA, COMPTEL and EGRET for
3C 279: $\alpha_X \approx 0.6$ with steepening at a few MeV, so that $\alpha_\gamma \approx 1.0$. On the other hand, a quasar like 3C 273 harbouring a “miniblazar”, but otherwise emitting predominantly thermal UV and IR emission, is very likely to show signs of interactions of these photons with the accelerated particles inside the nuclear jet. The interactions reduce the maximum energies of the particles they can obtain by diffusive acceleration limited by energy losses. As shown, this reduction moves the injection energy of secondaries at the top of the cascade closer to the energy where the jet becomes optically thick with respect to pair creation. Therefore 3C 273 develops a different cascade spectrum with two bumps at a few MeV and TeV, respectively. The latter is absorbed by infrared photons from dust clouds within the host galaxy, whereas the MeV bump has power law wings on both sides in energy space matching the observed X–ray and gamma ray power–laws consistent with GINGA, COMPTEL and EGRET observations. Of course, this does not rule out the existence of additional emission components in the X–ray band.

For 3C 273 flattening of the spectrum at a few GeV is expected, but observations do not seem to indicate this – a detail deserving further measurements. It is at the present stage of investigation unclear, what the variability properties should be like. Variability must be very complex, because we are actually dealing with a source extending over scales from fractions of a parsec to parsecs (transverse size of the jet or downstream emission regions of shocks). However, the variability time scale of the submm–nearinfrared target photons $t_{ir} \simeq r_{b,\perp}(1+z)/D_jc \simeq 180$ hours) should at least be contained in the X– and gamma ray autocorrelation functions, the amplitude should be significantly larger because of cascade reprocessing. An outburst in the power of the jet itself should manifest itself first by a flaring submm/infrared/optical continuum (the target) followed very shortly after by the gamma rays and then much later by a corresponding rise of the radio flux. Since the radio flux peaks at its self–absorption frequency, the delay at a given radio frequency determines the scale length of the radio emitting part of the jet, viz. $\Delta t \simeq (r - r_b)(1+z)/D_jc$ with $r = 4(\nu_{obs}/\nu_b)^{-1}$ pc and $\nu_b = 2 \cdot 10^{10}$ Hz.

As a corollary of the interpretation of the gamma rays from 3C 273 being of hadronic
origin it follows that 3C 273 should also be a source, which provides roughly 1% of a) the flux of cosmic rays at $10^{19}$ eV at Earth and b) the atmospheric background neutrino flux at 10 TeV in a solid angle of $1^\circ \times 1^\circ$. To come to these conclusions consider the following:

(i) In contrast to the electromagnetic flux the neutron as well as the neutrino flux density spectrum from photomeson production are flat (the luminosity increases proportional with the energy), because the UHE protons in the jet scatter on a polychromatic target, i.e. an inverse power law, so that the number of target photons increases rapidly with increasing energy. The electromagnetic cascade washes out the injection energy of the photons and redistributes the power below the energy above which the jet is optically thick with respect to pair creation.

(ii) Hadronic proton–photon interactions produce neutrons by isospin flip. The apparent neutron luminosity $L_n \approx L_{\pi^+}$ is given by $L_n \approx \frac{4}{13} L_\gamma$ (Ref. [13] where pion and pair production are taken into account) and the maximum observed energy is $E_n = \frac{1}{2} D_j E_{p,\text{max}}/(1+z)$. Hence one obtains using $\int_\nu F_\nu d\nu = 3 \cdot 10^{13}$ Jy Hz

$$F_{n\rightarrow p} = \frac{4}{13} E_{n,\text{max}}^{-1} \int_\nu F_\nu d\nu \simeq 7 \cdot 10^{-19} \text{ cm}^{-2} \text{ s}^{-1}$$

while the observed flux at $10^{19}$ eV is roughly $6 \cdot 10^{-17}$ cm$^{-2}$ s$^{-1}$ [16]. The neutrons are not confined to the magnetic field in the jet and because of the $\beta$–decay length $l_\beta = 100[\gamma_n/10^{10}]$ kpc they can escape the host galaxy without adiabatic losses (cf. original papers in Ref. [16]). Energy losses on the cosmic microwave background further reduce the flux above $10^{19}$ eV. Moreover, diffusion of the protons due to random magnetic fields in the surroundings of 3C 273 could smear the cosmic ray beam, thus reducing the cosmic ray flux by a factor $D_j^{-3} \simeq 3 \cdot 10^{-3}$ and the particle energy by $D_j^{-1} = 1/7$. Such smearing would, however, increase the total number of radio sources contributing to the entire extragalactic flux, since then also non–beamed (not core dominated, steep spectrum) sources would contribute (cf. Ref. [17] for a detailed model with direct proton escape from isotropic hot spots at the tip of powerful jets).

(iii) The neutrino luminosity for all species is given by $L_\nu \approx \frac{3}{13} L_\gamma$ and $E_{\nu,\text{max}} =
\( \frac{1}{20} D_f E_{p,\text{max}}/(1 + z) \), see Ref. [13]. This yields the flux

\[
F_\nu(p\gamma) = \frac{3}{13} E_{\nu,\text{max}}^{-1} \int_x^\gamma F_\nu d\nu \approx 5 \cdot 10^{-18} \text{cm}^{-2}\text{s}^{-1}
\]

important at Fly’s Eye energies [18]. The neutrino flux from \( pp \)-interactions (which dominates at energies in the 10 TeV range, see Ref. [9]) is given by

\[
F_\nu(pp) \approx \frac{10 r_g}{r_b \ln E_{p,\text{max}}/E_{p,\text{min}}} F_\nu(p\gamma) \left( \frac{E}{E_{\nu,\text{max}}} \right)^{-1} \approx 2 \cdot 10^{-5} m_9 F_\nu(p\gamma) \left( \frac{E}{10^{10} \text{GeV}} \right)^{-1}
\]

reaching \( 10^{-16} m_9 \text{ cm}^{-2}\text{s}^{-1} \) at 10 TeV. Thus, the neutrino flux from the jet of 3C 273 is far below the background neutrino flux from AGN as predicted by several authors (see Refs. [19,20]). However, the effective intensity for a solid angle of \( 1^\circ \times 1^\circ \) reaches \( 3 \cdot 10^{-17} m_9 \text{ cm}^{-2}\text{s}^{-1} \text{ sr}^{-1} \text{ GeV}^{-1} \) at 10 TeV for 3C 273 alone, compared to \( \approx 10^{-14} \text{ cm}^{-2}\text{s}^{-1} \text{ sr}^{-1} \text{ GeV}^{-1} \) for the hypothetical AGN background (which itself becomes comparable with the atmospheric background at 10 TeV). Considering that the number of similar flat–spectrum radio sources is very large [21], detection of the entire population of these powerful cosmic accelerators seems possible by refined techniques using horizontal atmospheric showers [22] and by taking advantage from the fact that the candidate source positions are known to any desired accuracy.

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FIGURES

FIG. 1. The broadband spectrum of 3C 273 from submm wavelengths up to hard gamma rays decomposed into a “blazar” component from a nuclear jet and thermal emission from accreted gas. Data for the submm-optical blazar component are taken from Ref. [3], who suggested the decomposition based upon polarization and variability, for the X–ray range from Ref. [23], for the hard X–ray/soft gamma ray range from Ref. [24] and, finally, for the gamma ray range from Ref. [1]. Several attempts to detect 3C 273 with Cherénkov telescopes have failed [25]. The synthetic spectrum shown comprises the emission from shock accelerated electrons (submm–optical) and protons (X–gamma rays). The emission mechanism is optically thin synchrotron emission from “primary” (=shock accelerated) electrons and from pairs produced by an unsaturated synchrotron cascade initiated by the protons cooling in the soft synchrotron photon soup. The TeV emission is absorbed by infrared photons from dust within the host galaxy. The model predicts a flattening of the spectrum between 10 GeV and a few hundred GeV.
TABLES

TABLE I. Physical conditions in the compact jet of 3C 273. Basic parameters for the calculation of physical conditions are: redshift \( z \), Lorentz factor of jet \( \gamma_j \), angle to the l.o.s. \( \theta \), proton/electron energy ratio \( \eta = L_\gamma/L_{ir/o} \), equipartition constant \( k_e \approx u_{rel}/u_B \), luminosity of jet in m.f. and rel. particles \( L_{44} \) in units of \( 10^{44} \) erg/s, observed opening angle of jet \( \Phi_{ob} \) and \( \nu_c \) comoving frame cutoff frequency of the synchrotron spectrum from the primary electrons. Inferred quantities (consistent with observations) are the observed apparent speed \( \beta_{ob} \), the Doppler factor of the jet \( D_j = [\gamma_j(1 - \beta_j \cos \theta)]^{-1} \), the primary electron break frequency \( \nu_b \), the photon/magnetic energy ratio \( a_s \) due to the radiation produced in the jet and due to an external luminosity of \( 10^{47} \) erg/s, resp., the projected length of the radiative jet \( r_{b,ob} \) where most of the emission with \( \nu > \nu_b \) comes from, the deprojected length \( r_b \), the (unresolved) transverse size of the jet \( r_\perp \), the target photon compactness in the comoving frame \( l \), the magnetic field \( B_b \) at \( r_b \), the electron and proton break Lorentz factors \( \gamma_{e,b} \) and \( \gamma_{p,b} \) and finally the observed photon energy \( E_{\gamma}^* \) where the jet becomes optically thick with respect to pair creation.

| Basic parameters | 3C 273 (miniblazar) | Comment, reference |
|------------------|---------------------|-------------------|
| \( z \)          | 0.158               | [21]              |
| \( \gamma_j \)   | (\( \geq \) 7)     | [2]               |
| \( \theta \) [deg]| (\( \geq \) 8)     | [2]               |
| \( \eta \)       | 15                  | \( L_\gamma/L_{ir} \) (this work) |
| \( k_e \)        | 1                   | \( u_{rel} \approx u_B \) assumed (this work) |
| \( L_{44} \)     | (\( \geq \) 50)    | [26] - hot spot A |
| \( \Phi_{ob} \) [deg]| 4                   | [2]               |
| \( \nu_c \) [Hz] | \( 10^{14} \)      | [3], cf. [5]     |

| Inferred quantities | (\( \Lambda_e = 5, \Delta = 5 \)) | Equation |
|---------------------|----------------------------------|----------|
| \( \beta_{ob} \)   | 7                                | [4] Eq. (1)          |
| \( D_j \)          | 7                                | [4] Eq. (6)          |
| \( \nu_b \) [Hz]   | \( 2 \cdot 10^{10} \)          | [4] Eq. (30)         |
| Parameter          | Value     | Reference |
|--------------------|-----------|-----------|
| $a_{s,\text{in}}$ | 0.004     | Eq. (24), intrinsic |
| $a_{s,\text{ex}}$ | 0.125     |           |
| $r_{b,\text{ob}}$ [pc] | 0.6       | Eq. (31) |
| $r_b$ [pc]        | 4         |           |
| $r_{b,\perp}$ [pc] | 0.04      |           |
| $l$                | $2 \cdot 10^{-6}$ | Eq. (18) |
| $B_b$ [G]         | 0.3       | Eq. (24) |
| $\gamma_{e,b}$    | $1.5 \cdot 10^2$ | Eq. (7) |
| $\gamma_{p,b}$    | $3 \cdot 10^{10}$ | Eq. (15) |
| $E_\gamma^*$(obs) [TeV] | 18       | Eq. (4) |