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THE ORIGIN OF COSMIC RAYS ABOVE $10^{18.5}$ eV

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ABSTRACT

We discuss the origin of ultra–high-energy cosmic rays (UHECRs) in the energy range above $10^{18.5}$ eV where it is indicated that the spectrum becomes protonic and extends smoothly up to at least $10^{19.5}$ eV and is consistent with a smooth extension to $10^{20.5}$ eV. The acceleration of the $10^{19.5}$ eV component must occur within ~1 Gpc. We rule out the production and escape of protons and neutrons from active galactic nuclei. Composition arguments make unlikely any origin in metal-rich environments such as rich clusters and the inner regions of galaxies. We dismiss the canonical extended halo models since such halos are almost never observed although diffuse halos have been seen in QSO absorption-line studies of metallic absorption lines. Large-scale shocks from explosions and winds are analyzed including those originating both recently and at earlier cosmological epochs. Large-scale shocks can work well only if they occur in microgauss fields. Hot spots and cocoons of radio sources are a plausible source for UHECRs with the principal uncertainty in both cases being the adopted or inferred magnetic field strength. Cosmic shocks are formed as structure develops during gravitational collapse of primordial perturbations such as are found in the standard cosmological models in, for example, both pancakes and collisions of hierarchical merging subunits. Cosmic shocks can be good sites for UHECR acceleration if there is a primordial field of order $\gtrsim 10^{-9}$ G or, again, if microgauss fields can be self-generated in shocks. Table 1 summarizes the results of our analysis of all the conventional possibilities and indicates that jets, radio source cocoons, structure formation in clusters superclusters, and large-scale structures are all reasonable sources for the production of UHECRs.

We examine the possibility of a second-stage diffuse acceleration process from an ensemble of shocks that might explain the highest energy particles above 100 EeV by boosting their energy from 30 to 300 EeV into a diffuse isotropic component. Table 2 summarizes our results and indicates that at the present time there is as yet no such candidate ensemble of shocks that has both the power and the volume-filling factor to achieve this effect. The highest energy cosmic rays must therefore have been accelerated in one stage from their parent object.

Subject headings: acceleration of particles — cosmic rays — shock waves

1. INTRODUCTION: THE COMPOSITION METALLICITY CORRELATION

The discovery that the composition of ultra–high-energy cosmic rays (UHECRs) changes above $10^{18.5}$ eV to a proton-dominated component (Bird et al. 1993) and smoothly continues to at least $10^{19.5}$ eV and plausibly to $10^{20.5}$ eV allows significant insight into the origin of these particles. Previous studies of the acceleration of the UHECRs that have concentrated on the physics of various extragalactic shocks have had some predictive success concerning the nature and composition of the UHECRs (Cavallo 1978; Axford 1991, 1992, 1994; Biermann 1993a, b; Rachen & Biermann 1993; Rachen, Stanev, & Biermann 1993; Stanev, Biermann, & Gaisser 1993). An alternative suite of models has focused on acceleration processes that occur in the nuclei of active galaxies where the high-energy particles that escape are neutrons, which subsequently decay into protons (Biermann & Strittmatter 1987; Sikora et al. 1987; Begelman, Rudak, & Sikora 1990; Protheroe & Szabo 1992; Szabo & Protheroe 1994). Sommers (1994), Elbert & Sommers (1995) and Sigl, Schramm, & Battacharjee (1994) discussed both cases in relation to the production of the highest energy cosmic rays and some of the conclusions are similar to those reached here. In this paper we show that neither protons nor neutrons of the highest energy can escape from active galactic nuclei (AGNs) because they photodestruct on a central radiation field (§ 2). Moreover, in the central regions of AGNs, synchrotron-Compton losses and photopion losses restrict the maximum energy to which particles can be accelerated to well below the observed highest energies (§ 3). Thus the identification of the highest energy particles as protons places a severe constraint on the shocks at which the acceleration can occur.

The identification of the UHECRs as protons also raises the question as to why they are not heavy nuclei. The suggestion that UHECRs should be heavier nuclei is based on two simple arguments: the composition of shock-accelerated ions reflects the composition of the ambient medium, and the high-energy cutoff for shock acceleration increases with the charge number, Z, of the nucleus. Here we argue that the observed composition of UHECRs implies that they are accelerated in regions of very low metallicity, so that the absence of heavy nuclei in the UHECRs is attributed to the absence of heavy ions in the ambient plasma. The alternative possibility that the UHECRs are accelerated as heavy nuclei and then photodissociate into protons (and neutrons that decay into protons) is excluded because the energy loss rates for protons and the photodissociation rate for heavy nuclei ^{56}Fe are comparable at the energies of $10^{19.5}$ eV (Puget, Stecker, & Bredekamp 1976, Fig. 2 and Fig. 8; Berezinskii et al. 1990, § 4.4c and Fig. 4.10; Elbert & Sommers 1995). Consequently, for the cosmic-ray spectrum

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extending smoothly above and below $10^{19.5}$ eV there is a strong correlation between metallicity at the source and the observed composition of cosmic rays.

The composition argument then becomes quite powerful and one can immediately rule out more extended extranuclear acceleration sites that have any significant metallicity. This point eliminates as the site of the UHECRs the near environments of galaxies including our galaxy, starburst galaxies, the inner (but nonnuclear) regions of active galaxies and radio galaxies, rich clusters of galaxies, and radio sources in these clusters. The interesting suggestion that colliding and interacting galaxies may be the source of UHECRs (Cesarsky & Ptuskin 1993; Cesarsky 1994) is most likely ruled out by the same arguments since the shocks occur in metal-rich material unless they are at very large distances ≥ 50 kpc from the center of the system. We know that metals are present at ≤ 50 kpc from QSO absorption line studies of C IV and Mg II (Steidel 1995; Bergeron 1995).

The fact that the distribution of protons extends smoothly up to $10^{19.5}$ eV places limits on the cosmic volume in which these particles can be accelerated. A proton of this energy cannot travel more than ~ 1 Gpc (Berezinskii et al. 1990, Fig. 4.9) due to pair production and/or pion production on the microwave background and a similar limit is placed on ^{56}Fe at the same energy (Puget et al. 1976; Elbert & Sommers 1995). This places the sources at redshifts $z \leq 0.3\ h$ ($h = H_0/100$ km s^{-1} Mpc^{-1}) and therefore significant cosmological evolution can be ignored.

We consider the acceleration of particles near large-scale shocks. Due to the very large gyroradii of the UHECRs, sites for their acceleration must either be very large or have very strong magnetic fields, to ensure that the inevitable expansion losses in the postshock flow are compensated by diffusive shock acceleration (see § 4) at the shock front. One possibility for very large sources is extended galactic halos (e.g., Jokipii & Morfill 1992). However, despite intensive searches there is no evidence for a ubiquitous extended halo component around normal galaxies. The few observed halos such as NGC 891 and NGC 4631 (see Wang et al. 1995) are regarded as special cases (Norman & Ikeuchi 1989). The most reasonable assumption is that in general extended halos do not exist. What are seen are thick disks (Norman 1995). Smaller scale halos are ruled out as a source of UHECRs by the metallicity argument. In addition, Axford (1994 plus references) has ruled out the Jokipii & Morfill (1992) suggestion of extended galactic halos as a production site of UHECRs on energetic grounds.

We are left then with the task of establishing extended extranuclear sources for UHECRs in regions of low or zero metallicity at redshifts less than a few tenths. We examine large-scale shocks generated by explosions or winds (§ 4). We review and refine the canonical radio galaxy hot spot model and show that overpressured cocoons can also be sources of UHECRs (§ 5). We analyze the acceleration of UHECRs in shocks created in the collapse and formation of clusters and superclusters (§ 6).

Our general procedure is to estimate whether UHECR energies can be attained as the maximum energy of an accelerated particle. This is limited by expansion losses or the finite size (and age) of the accelerating shock. We then calculate further constraints on the details of the acceleration process.

2. UHECRs CANNOT ESCAPE AGNs

We consider the model (Protheroe & Szabo 1992; Szabo & Protheroe 1994) in which UHECR acceleration occurs in AGNs where the large photon flux allows significant conversion into neutrons in photopion production, with the neutrons escaping, ultimately being reconverted into protons through β-decay. We show that it is unlikely that any UHECRs above 10^{19} eV can escape.

A reasonable limit on the maximum energy for accelerated particles in AGNs as a function of the luminosity L can be calculated as follows. Assuming equipartition between the energy density in radiation and the magnetic field energy density (e.g., Rees 1987) in a source of radius R implies

$$\frac{L}{4\pi R^2 c} = \frac{B^2}{8\pi}\,. \tag{1}$$

The maximum energy achievable due to diffusive shock acceleration in a shock with velocity $V_s = \beta_s c$ and size R_s with saturated turbulence providing the scattering (with the particle mean-free path equal to its gyroradius) is

$$E_{\rm max} \sim Ze\beta_s BR_s\,. \tag{2}$$

This limiting energy results from either the finite size (or age) of the shock or due to expansion losses in the postshock flow, but neglects radiation losses. Assuming $R_s \simeq R$, this can be reformulated using equation (1) in equation (2) as

$$E_{\rm max} = Ze\beta_s\left(\frac{2L}{c}\right)^{1/2}, \tag{3}$$

which gives

$$E_{\rm max} = 2.5 \times 10^{19}\beta_{-1} ZL_{46}^{1/2}\ {\rm eV}\,, \tag{4}$$

with $L_{46} = L/10^{46}$ ergs s^{-1} and $\beta_{-1} = \beta_s/0.1$. If one interprets the proton-dominated component from $10^{18.5}$ to $10^{19.5}$ eV in the Fly's Eye data (Bird et al. 1993) as coming from AGNs, then only fast shocks in luminous AGNs, with $L \geq 10^{44}$–10^{46} ergs s^{-1} and $\beta_s \geq 0.1$, can account for this component.

In the hard photon spectra associated with AGNs, escaping hadrons are degraded by photopion production. The optical depth for photopion production on the intense radiation field (assuming that νF_ν is approximately constant) is

$$\tau = \frac{\sigma_\pi L}{4\pi R\langle\epsilon\rangle c}\,, \tag{5}$$

where $\langle\epsilon\rangle$ is the mean photon energy and $\sigma_\pi \simeq 0.1$ mb is the cross section for photopion production.

The compactness parameter, l, is defined for AGNs as

$$l = \frac{\sigma_{\rm T} L}{4\pi m_e c^3 R}\,, \tag{6}$$

where $\sigma_{\rm T}$ is the Thomson cross section. Then equation (5) gives

$$\tau = 760 l\left(\frac{\langle\epsilon\rangle}{0.1\ {\rm eV}}\right)^{-1}. \tag{7}$$

The threshold for photopion production on infrared photons is (Biermann & Strittmatter 1987)

$$E_{\rm tr} \simeq 10^{16}\left(\frac{\lambda}{1\ \mu{\rm m}}\right){\rm eV}\,. \tag{8}$$

For most AGNs one infers a compactness parameter $l \geq 1$ (e.g., Done & Fabian 1989) in the central region where the initial acceleration occurs, so neutrons above $\simeq 10^6$–10^7 GeV cannot escape. This agrees with Szabo & Protheroe (1994).

Even if the hadrons were to escape the region of compactness parameter of order unity, then they would certainly be stopped by the radiation field in the broad line region (BLR) and the radiation field from the obscuring disk or torus. Previous analyses have assumed that outside the central region the photon trajectories are almost radial. This ignores the reprocessed component of the radiation field that is absorbed and reradiated by the clouds that constitute the BLR. The canonical covering factor for these clouds is 10%. There is also a component of the radiation field that is absorbed by the warped disk or obscuring torus and then reradiated (see Pier & Krolik 1993).

If we assume conservatively that a fraction, $\eta \sim 10\%$, of the continuum is absorbed and then reemitted at 1–10 μm, depending on how close the absorbers are to the central source, then the optical depth for photodestruction of hadrons in this region on the isotropic component is

$$\tau_i = \frac{\eta \sigma_\pi L}{4\pi R c \langle \epsilon_i \rangle}\,, \tag{9}$$

where subscript i denotes the isotropic component. This gives

$$\tau_i = 0.6 \eta_{-1} L_{46} R_{\rm pc}^{-1} \left(\frac{\langle \epsilon_i \rangle}{0.1\ {\rm eV}}\right)^{-1}\,, \tag{10}$$

$\eta_{-1} = \eta/0.1$ and $R_{\rm pc} = R/1$ pc. Thus the UHECR component does not escape unless the source luminosities are lower than 10^{46} ergs s^{-1}. In either case, the spectrum of escaping protons would not extend smoothly up to $10^{19.5}$ eV as observed in the Fly's Eye spectrum.

3. LOSSES AND THE MAXIMUM ENERGY

The maximum energy to which cosmic rays can be accelerated in the absence of losses is given by equation (2). Close to the black hole powering AGNs, losses can limit the maximum energy to a much lower value.

The maximum energy attainable is determined either by a balance between the energy gain due to shock acceleration and the energy loss incurred between subsequent shock crossings, or, if losses are absent, by the time available for acceleration. For relativistic particles with scattering mean free path $\lambda_f \simeq R_g$, where $R_g \sim E/ZeB$ is the gyroradius, the standard theory of shock acceleration (e.g., Drury 1983; Blandford & Eichler 1987) yields the net energy gain (up to factors of order unity) due to these combined effects:

$$\frac{dE}{dt} = \frac{Ze}{c} B V_s^2 - \epsilon \frac{V_s}{R_s} E - \dot{E}_{\rm loss}\,. \tag{11}$$

The first term on the right-hand side of equation (11) corresponds to diffusive shock acceleration at a shock with velocity V_s, the second term describes expansion losses in the postshock flow, which cannot be avoided in a propagating, nonplanar shock. The value of the parameter $\epsilon \leq 1$ depends on the detailed conditions near the shock. The last term describes energy losses in the intense radiation field or magnetic fields in AGNs.

Synchrotron-Compton losses of nuclei with charge Ze and mass Am_p correspond to an energy loss rate

$$\dot{E}_{\rm loss} = \frac{1}{t_{\rm rad}} \left(\frac{E^2}{mc^2}\right)\,, \tag{12}$$

where

$$t_{\rm rad} = \left(\frac{m_p}{m_e}\right)^3 \frac{A^3}{Z^4} \left(\frac{6\pi m_e c}{\sigma_{\rm T} B^2}\right) \tag{13}$$

is the radiation time. The energy loss time of a particle with Lorentz factor $\gamma = E/mc^2$ is $t_{\rm loss} = t_{\rm rad}/\gamma$.

If one balances the shock acceleration rate with synchrotron-Compton losses one finds a maximum energy

$$E_{\rm sc} \sim \sqrt{mc^2}(ZeB\beta_s V_s t_{\rm rad})^{1/2}\,. \tag{14}$$

By combining equation (2), which applies in the absence of losses, and equation (14) one finds

$$E_{\rm sc} \sim (\sqrt{mc^2 E_{\rm max}}) \left(\frac{V_s t_{\rm rad}}{R_s}\right)^{1/2}\,. \tag{15}$$

In applying this to particle acceleration in AGNs we use equation (1) to estimate the magnetic field in equation (13). Also we use equation (6) to eliminate the size R of the source region from the estimates, on the grounds that variability studies allow an estimate of the compactness parameter (e.g., Done & Fabian 1989). This gives

$$R_{\rm pc} = 7 \times 10^{-3}(L_{46}/l)\,. \tag{16}$$

The typical equivalent magnetic field is

$$B \sim 37(l/L_{46}^{1/2})\ {\rm G}\,, \tag{17}$$

and the radiation time is

$$t_{\rm rad} = \frac{3A^3}{4Z^4} \left(\frac{m_p}{m_e}\right)^3 \left(\frac{R}{lc}\right) \sim 10^8 \frac{A^3}{Z^4} \left(\frac{L_{46}}{l^2}\right) {\rm yr}\,, \tag{18}$$

corresponding to a loss time for particles with $E \sim E_{\rm max}$ of the order of days:

$$t_{\rm max} = \left(\frac{mc^2}{E_{\rm max}}\right) t_{\rm rad} \sim 4 \times 10^{-2} \frac{A^4}{Z^5} \left(\frac{L_{46}^{1/2}}{\beta_{-1} l^2}\right) {\rm yr}\,. \tag{19}$$

Using equations (15) and (19) one finds the limiting energy due to synchrotron-Compton losses for $R_s \sim R$

$$E_{\rm sc} \sim E_{\rm max} \left(\frac{V_s t_{\rm max}}{R_s}\right)^{1/2}\,, \tag{20}$$

which is

$$E_{\rm sc} \sim 3.5 \times 10^{18} \beta_{-1} \frac{A^2}{Z^{3/2}} \left(\frac{L_{46}^{1/4}}{l^{1/2}}\right) {\rm eV}\,. \tag{21}$$

The synchrotron-Compton losses limit the maximum energy to less than the value allowed in the absence of losses when $E_{\rm sc}$, as given by equation (21), is less than $E_{\rm max}$, as given by equation (2). This is the case when $V_s t_{\rm max} < R_s$, and the compactness parameter in the source region satisfies

$$l \geq 1.8 \times 10^{-2} \frac{A^4}{Z^5} L_{46}^{-1/2}\,, \tag{22}$$

independent of shock speed. Expressed in terms of the radius (see eq. [16]), this corresponds to

$$R \leq 0.38 \frac{Z^5}{A^4} L_{46}^{3/2}\ {\rm pc}\,. \tag{23}$$

The parameter that ultimately determines the relative importance of synchrotron-Compton losses is

$$\frac{V_s t_{\rm rad}}{R_s} \sim 4.8 \times 10^8 \beta_{-1} \frac{A^3}{Z^4} l^{-1}\,, \tag{24}$$

which depends only on the compactness parameter and shock speed.

We also need to consider losses due to photopion production. This process has a loss length l_π which can be expressed in the pion production cross section $\sigma_{\pi\gamma}$ and the inelasticity $K_{\pi\gamma} = \| \Delta E \|/E$ of the reaction as

$$l_\pi = \frac{1}{n_{\rm ph}\langle \sigma_{\pi\gamma} K_{\pi\gamma}\rangle} = 2.8 \times 10^{13}\left(\frac{\langle\epsilon\rangle}{0.1 \text{ eV}}\right)\left(\frac{L_{46}}{l^2}\right) \text{ cm} . \quad (25)$$

where equations (6) and (5) are used and the photon density is estimated as $n_{\rm ph} \sim L/4\pi R^2 c\langle\epsilon\rangle$. The corresponding loss rate is

$$\dot{E}_{\rm loss} = \left(\frac{dE}{dt}\right)_\pi \sim -\frac{c}{l_\pi} E . \quad (26)$$

Balancing the photopion losses with the acceleration rate at the shock, one finds the corresponding maximum energy

$$E_\pi \sim ZeB\beta_s^2 l_\pi . \quad (27)$$

Defining the optical depth of the shock for pion production as $\tau_{\pi s} \equiv R_s/l_\pi$ one can write this in terms of $E_{\max}$ (see [2]), as

$$E_\pi \sim E_{\max}\left(\frac{\beta_s}{\tau_{\pi s}}\right) . \quad (28)$$

Thus, losses due to pion production limit the maximum energy a cosmic ray can attain to $E_\pi < E_{\max}$ whenever

$$\tau_{\pi s} \geq \beta_s , \quad (29)$$

which is a rather stringent condition. Using equation (5) with $R \sim R_s$ gives

$$l \geq 1.3 \times 10^{-4}\beta_{-1}\left(\frac{\langle\epsilon\rangle}{0.1 \text{ eV}}\right) . \quad (30)$$

The estimate (27) corresponds to

$$E_\pi \sim 3.3 \times 10^{15}\beta_{-1}^2\left(\frac{L_{46}^{1/2}}{l}\right)\left(\frac{\langle\epsilon\rangle}{0.1 \text{ eV}}\right) \text{ eV} . \quad (31)$$

For protons ($Z = A = 1$) losses due to pion production are more important than synchrotron-Compton losses if $E_\pi < E_{\rm sc}$, or equivalently if

$$l \geq 10^{-6}\beta_{-1}^2 L_{46}^{1/2}\left(\frac{\langle\epsilon\rangle}{0.1 \text{ eV}}\right)^2 . \quad (32)$$

We conclude that the effect of losses near the black hole effectively rules out the acceleration of UHECRs.

4. EXPLOSIONS AND WINDS

In estimating the acceleration of UHECRs resulting from explosions and winds associated with quasars and AGNs we assume spherical symmetry, and use the appropriate similarity solutions $R_s \sim t^\alpha$. The shock waves are assumed to propagate through the intergalactic medium (IGM) with density ρ and with a large-scale magnetic field of strength B. We subsequently refine these assumptions where necessary.

We express the energy gain in terms of an energy gain per unit shock radius rather than time, using $d/dt = V_s(d/dR_s)$, and neglect radiation losses. Equation (11) then reduces to

$$\frac{dE}{dR_s} \approx \frac{ZeBV_s}{c} - \epsilon\frac{E}{R_s} . \quad (33)$$

If the shock propagates into a stationary medium, standard shock acceleration theory (e.g., Drury 1983) gives an estimate for the parameter ϵ which detemines the relative importance of expansion losses:

$$\epsilon \sim \frac{t_+}{t_+ + t_-} = \frac{r}{r + r_B} , \quad (34)$$

where $r = \rho_+/\rho$ and $r_B = B_+/B$ are the density and field compression ratios at the shock. The quantities t_- (t_+) correspond to the times spent by the particle in the upstream (downstream) flow during each cycle of two subsequent shock crossings. Without additional amplification of the postshock magnetic field, one has $r_B \leq r$ and $\epsilon \geq 1/2$.

For a simple scaling law $B \propto R_s^{-\nu}$ a similarity solution has $BV_s \propto R_s^{-\mu}$ with $\mu = \nu + 1/\alpha - 1$. Direct integration of equation (33) then yields

$$E(R_s) = E_i \mathscr{R}^{-\epsilon} + \left(\frac{Ze}{c}\right)V_s BR_s \psi(\mathscr{R}) , \quad (35)$$

where E_i is the injection energy, and $\mathscr{R} \equiv R_s/R_i \geq 1$ is the expansion factor of the blast wave since injection. The function $\psi(\mathscr{R})$ is defined as

$$\psi(\mathscr{R}) = \begin{cases} (1 - \mathscr{R}^{-\chi})/\chi , & \text{when } \chi \neq 0 , \\ \ln \mathscr{R} , & \text{when } \chi = 0 , \end{cases} \quad (36)$$

with $\chi = 1 + \epsilon - \mu$. The quantity $E(R_s)$ is interpreted as the energy where the distribution of shock-accelerated particles terminates when the blast wave has expanded to radius R_s. We neglect the decaying term involving the injection energy E_i.

For a spherical Sedov-Taylor blast wave with explosion energy E, one has $R_s \sim (E/\rho)^{1/5}t^{2/5}$. Taking $\epsilon = 1/2$ and $\nu = 0$, one has $\chi = 0$ and

$$E_{\max} = \left(\frac{Ze}{c}\right)B\left(\frac{E}{\rho}\right)^{2/5} t^{-1/5} \ln \mathscr{R} , \quad (37)$$

which gives

$$E_{\max} = 5 \times 10^{17} ZB_{-8} E_{61}^{2/5} \rho_{-31}^{-2/5} t_8^{-1/5} \ln \mathscr{R} \text{ eV} . \quad (38)$$

Here subscripts give normalizing values in cgs units except for time which is measured in years.

Clearly, unless the field strength is of order ~ 1 μG, these explosions cannot produce UHECRs. In rich clusters such field strengths are observed (e.g., Dreher, Carilli, & Perley 1987; Taylor et al. 1990). However, rich clusters are eliminated as sources for UHECRs by the metallicity argument, as metallicities of order solar are now observed in such clusters (Fabian 1988).

A similar situation occurs for winds. If the mechanical luminosity driving the wind is L_w a spherical wind bubble expands as $R_s \sim (L_w/\rho)^{1/5}t^{3/5}$. Taking $\epsilon \approx 1/2$ and $\nu = 0$ one has $\chi \approx 5/6$ and $E_{\max}$ is then

$$E_{\max} = 1.2\left(\frac{Ze}{c}\right)B\left(\frac{L_w}{\rho}\right)^{2/5} t^{1/5} , \quad (39)$$

which is

$$E_{\max} = 10^{18} ZB_{-8} L_{46}^{2/5} \rho_{-31}^{-2/5} t_8^{1/5} \text{ eV} . \quad (40)$$

For remnant shock waves of explosions from earlier epochs such as the quasar phase at $z = 2$, a similar situation occurs. For an $\Omega = 1$ cold dark matter dominated universe with a baryonic content $\Omega_g \ll 1$, the expansion law is dominated by the dynamics of the cold dark matter shell expanding into a density decreasing as $\rho = (6\pi Gt^2)^{-1}$ and a magnetic field decaying as $B \propto t^{-4/3}$ due to the Hubble expansion. Such remnant bubbles expand as $R_s \sim (6\pi GE)^{1/5} t^{4/5}$ (Bertschinger 1983; Ostriker & McKee 1988; Voit 1994). The corresponding estimate of $E_{\rm max}$ still follows from an expression of the form (14). However, to take account of the expansion losses in the preshock Hubble flow and the field decay $B \propto R_s^{-5/3}$ one has to redefine the parameter χ as $\chi = -(1 + 2\epsilon)/12 < 0$. This means that in this case, the maximum energy is set by the time of injection t_i. For $\epsilon \sim 1/2$ one has $\chi = -1/6$ and

$$E_{\rm max} = 6\left(\frac{Ze}{c}\right) B_i (6\pi GE)^{2/5} t_i^{3/5} , \quad (41)$$

where $B_i \equiv B(t_i)$. This becomes

$$E_{\rm max} = 1.3 \times 10^{18} Z B_{-8} E_{62}^{4/5} t_{10}^{3/5} \text{ eV} , \quad (42)$$

where $t_{10} \equiv t_i/(10^{10}$ yr) and $B_{-8} = B_i/(0.01\ \mu$G). However, it turns out that at late times the velocity of the shock relative to the Hubble flow is only ~ 100 km s^{-1} and it is unlikely that a strong accelerating shock can act on the IGM.

We conclude that these large-scale blast waves are not sufficient to be the source of the UHECRs. The cosmic magnetic field that they propagate into is uncertain. In the centers of clusters the fields can be in equipartition with the thermal gas and have field strengths $\geq 10\ \mu$G (Ge and Owens 1993). A problem with using these regions is that they are often of high metallicity so that the acceleration would have to occur during the collapse virialization and merging of subunits before significant enrichment occurred. It is probable that in many cases the postshock flow is turbulent, leading to the generation of a postshock magnetic field strength far above the value due to shock compression alone. In the case of extreme self-generated fields, so that effectively $r_B \gg r$, one has $\epsilon \sim r/r_B \ll 1$ and postshock expansion losses can be neglected. Particles undergoing diffusive shock acceleration then spend almost all of their time in the (undisturbed) medium upstream from the shock ($t_+ \ll t_-$). In this case equations (35) and (36) still apply with $\epsilon = 0$ and $\chi = 1 - \mu = 2 - \nu - 1/\alpha$. The maximum energy is now solely determined by the time of injection or the age of the shock. This conclusion does *not* apply to cosmological shock waves, where expansion losses in the upstream Hubble flow cannot be avoided, leading to $\epsilon \sim 1$ regardless the value of r_B.

The maximum energy in the case of a Sedov-Taylor blast wave ($\chi = -1/2$) and a continuously driven spherical wind ($\chi = 1/3$) in a constant upstream magnetic field ($\nu = 0$) for $\mathscr{R} = R_s/R_i \gg 1$ is then

$$E_{\rm max} \approx \begin{cases} 2\left(\frac{ZeB}{c}\right)\left(\frac{E}{\rho}\right)^{2/5} t_i^{-1/5} & \text{Sedov-Taylor} , \\ 3\left(\frac{ZeB}{c}\right)\left(\frac{L_w}{\rho}\right)^{2/5} t^{1/5} & \text{driven wind} . \end{cases} \quad (43)$$

This yields in these two cases

$$E_{\rm max} \approx 2.5 \times 10^{18} Z B_{-8} \rho_{-31}^{-2/5} \text{ eV} \times \begin{cases} E_{62}^{2/5} t_{i8}^{-1/5} , \\ L_{46}^{2/5} t_8^{1/5} . \end{cases} \quad (44)$$

Again, these blast waves seem incapable of generating the UHECRs except for magnetic field strengths in excess of a few microgauss.

5. RADIO HOT SPOTS AND COCOONS

There are three possible acceleration sites in a typical powerful F-R II type radio source—the hot spot capping the jet, the bow shock preceding this hot spot in the IGM, and the shock bounding the cocoon of waste material left by the jet. Hot spots with very strong fields can be inferred for those with optical synchrotron emission (Meisenheimer et al. 1989). Biermann and colleagues (e.g., Rachen & Biermann 1993) have pointed out their advantageous properties for accelerating UHECRs, although the high field strengths have been questioned by Axford (1994).

For Mach disks powering hot spots, and the bow shock preceding the hot spot in the IGM, the factor limiting the energy gain is the transverse size R_j of the shock. Particles drift along the shock surface until they reach the edge and escape (see Jokipii 1987). The associated maximum energy follows from equation (2):

$$E_{\rm max} \sim \frac{ZeBV_j R_j}{c} , \quad (45)$$

which corresponds to the potential drop across the face of a quasi-normal shock. For typical hot spot parameters one finds (Biermann & Strittmatter 1987)

$$E_{\rm max} = 10^{20} B_{-4}\left(\frac{R_j}{1 \text{ kpc}}\right)\beta_j \text{ eV} , \quad (46)$$

with $\beta_j = V_j/c$, where V_j is the velocity of the jet.

The hot spot of shocked jet material, and the bow shock preceding it advance through the IGM with a speed, V_h, given by

$$V_h \approx \left(\frac{L_j}{\pi R_{\rm bs}^2 \rho_{\rm IGM} V_j}\right)^{1/2} , \quad (47)$$

where L_j is the mechanical luminosity (thrust) of the jet. This relation follows from ram-pressure balance for a bow shock of size $R_{\rm bs}$ (Chakrabarti 1988). Using this velocity one finds

$$E_{\rm max} = 10^{18} Z L_{46}^{1/2} \beta_j^{-1/2} n_{-4}^{-1/2} B_{-6} \text{ eV} , \quad (48)$$

where n is the number density of the IGM. Note that equation (48) implies that $E_{\rm max}$ is independent of the size of the bow shock. We conclude that bow shocks cannot account for UHECRs.

The cocoon of waste material processed in the hot spots expands sideways with a radius, R_c, of order (Begelman & Cioffi 1989; Achterberg, Melrose, & Norman 1995)

$$R_c \sim \left(\frac{L_j}{\rho_{\rm IGM} V_h}\right)^{1/4} t^{1/2} , \quad (49)$$

at time t after the passage of the jet. The maximum energy of particles produced at the cocoon follows from equations (13) and (14) with $\chi = \epsilon - \nu$. On the one hand, if there is no self-generation of field in the cocoon one finds, for $\epsilon = \frac{1}{2}$ and $\nu = 0$,

$$E_{\rm max} = 2\left(\frac{Ze}{c}\right) B \left(\frac{L_j}{\rho_{\rm IGM}} V_h\right)^{1/2} . \quad (50)$$

Expressing the speed of advance as $V_h \sim D_j/t$ with D_j the length of the jet and t the age of the source, one finds

$$E_{\rm max} \approx 5 \times 10^{19} Z B_{-6} L_{46}^{1/2} t_8^{1/2} n_{-4}^{-1/2} D_{10}^{-1/2} \text{ eV} , \quad (51)$$

where $D_{10} \equiv D_j/(10$ kpc). On the other hand, assuming a fully turbulent cocoon with a strong self-generated magnetic field or field mixed in by turbulent diffusion from the processed jet material, implies $r_B \gg 1$ and $\epsilon \approx 0$. Then one finds

$$E_{\rm max} = \left(\frac{ZeB}{c}\right)\left(\frac{L_j}{\rho_{\rm IGM} V_h}\right)^{1/2} \ln \mathscr{R} , \quad (52)$$

which depends only logarithmically on time. This corresponds to

$$E_{\max} \sim 2 \times 10^{19} Z B_{-6} L_{46}^{1/2} t_8^{1/2} n_{-4}^{-1/2} D_{10}^{-1/2} \ln \mathscr{R} \text{ eV} . \quad (53)$$

In some cases the magnetic fields inferred for the shells of the cocoons greatly exceed 1 μG (Carilli, Perley, & Dreher 1988; Taylor et al. 1990; Taylor & Perley 1994).

We conclude that both cocoons and hot spots are favorable sites for UHECR acceleration. The field strength is the crucial but most uncertain parameter that determines the maximum energy in both cases.

6. SHOCKS RESULTING FROM STRUCTURE FORMATION

As the gravitating structures in the universe form by the continuous process of hierarchical clustering large-scale shocks can form in the gaseous component as cooler merging subunits are shock-heated to the new virial temperature of the new higher binding energy system. For a good discussion of this process in the context of galaxy and cluster formation see White & Frenk (1991), Bond et al. (1991), and Lacy & Cole (1993) where there are good explanations of the process of structure formation, biasing, and the concept of hierarchical clustering.

For a nonlinear collapsing structure in an expanding universe with (excess) mass M and a scale R, the typical shock velocity is of the order of

$$V_s \simeq V_{\rm ff} = \sqrt{\frac{2GM}{R}} . \quad (54)$$

An estimate of the maximum energy of the accelerated cosmic rays in the structure-forming process is then

$$E_{\max} = \frac{ZeB_s}{c} (GMR_s)^{1/2} . \quad (55)$$

For standard biased cold dark matter (CDM) in the modified Press-Schechter formalism given in the above references (see also Peebles 1976, 1993, Ch. 19) and for a top-hat distribution we can estimate:

$$R_s \simeq \left(\frac{3M}{4\pi\rho\delta}\right)^{1/3} = \left(\frac{2GM}{\Omega\,\delta H^2}\right)^{1/3} , \quad (56)$$

where Ω is the cosmological density parameter, H is the Hubble constant, and δ is the overdensity parameter relative to the background after collapse and virialization. This yields

$$E_{\max} = \left(\frac{ZeB_s}{c}\right) G^{2/3} M^{2/3} \delta^{-1/6} \Omega^{-1/6} H^{-1/3} (1+z)^{-1/2} . \quad (57)$$

A good estimate with the appropriate normalization for the mass at the peak of the distribution of structures at redshift z is given by

$$M = 10^{15} \left[\frac{0.4}{b(1+z)}\right]^{6/(n+3)} M_\odot , \quad (58)$$

where $b \simeq 2$ is the bias parameter and n is the spectral index of the initial cosmological density fluctuations.

For the Harrison-Zeldovich spectrum, $n = -1$, the maximum energy achievable is

$$E_{\max} = 3 \times 10^{18} Z B_{-6} \delta^{-1/6} \Omega^{-1/6} h^{-1/3} (1+z)^{-5/2} \text{ eV} . \quad (59)$$

Here $h \equiv H_0/(100 \text{ km s}^{-1} \text{ Mpc}^{-1})$.

In addition, there are even larger scale structures in the universe that may have collapsed in one dimension. Cosmological models show considerable evidence for pancake structures. Even in models such as CDM and hot dark matter (HDM), collapsing structures occur over a wide range of scales (e.g., Kofman et al. 1992). Particles can be accelerated at the strong shocks that form at, for example, pancake caustic surfaces (e.g., Doroshkevich, Shandarin, & Saar 1978).

If the magnetic field in the shocked pancake is given by the flux-freezing condition

$$B = B_0 \left(\frac{R}{w}\right)^2 , \quad (60)$$

where w is the width of the shocked pancake. If we assume an intergalactic field of $B_0 \sim 10^{-9}$ G (Kronberg 1994 and choose for typical parameters for pancaking of large-scale structure at the present epoch (see Kofman et al. 1992 plus references) $w \sim 0.03R$ and $R \sim 10$ Mpc, we find $B \simeq 1$ μG. Since $E_{\max} \propto V_s R$, the largest structures potentially produce the most energetic particles. For (proto-)superclusters one has $V_s \approx 1000$ km s^{-1} and

$$E_{\max} = ZeB\beta_s R_s \simeq 2\,10^{19} Z B_{-6} R_{10} \text{ eV} . \quad (61)$$

Here $R_{10} = R/10$ Mpc.

At these energies, losses due to pair production and/or pion production on the cosmic microwave background start to become important (see Berezinskii et al. 1990, Fig. 4.9). Pair production dominates below 3×10^{19} eV and has an associated loss length $l_{\rm pair} \simeq 1$ Gpc in the range 3×10^{18}–3×10^{19} eV. Pair production losses limits the particle energy whenever (cf. eq. [27])

$$E_{\rm pair} \equiv ZeB\beta_s^2 l_{\rm pair} \simeq 10^{20} Z \beta_{-2}^2 B_{-6} \text{ eV} \quad (62)$$

lies in the above energy range and *provided* the shock size exceeds

$$R_s \geq \beta_s l_{\rm pair} \simeq 10\beta_{-2} \text{ Mpc} . \quad (63)$$

Here $\beta_{-2} = \beta_s/(10^{-2})$. Such shocks produce particles with $E \leq E_{\rm pair} < E_{\max}$.

Pion losses become important for $E \geq 3 \times 10^{19}$ eV. The associated loss length, $l(E)$, is given by

$$l(E) \sim l_\pi e^{E_{\rm tr}/E} , \quad (64)$$

where

$$l_\pi = \frac{1}{n_{\rm ph}\langle\sigma_{\pi\gamma} K_{\pi\gamma}\rangle} \sim 10(1+z)^{-3} \text{ Mpc} , \quad (65)$$

using the value for $\langle\sigma_{\pi\gamma} K_{\pi\gamma}\rangle$ of Figure 3 of Rachen & Biermann (1993). For protons, the threshold energy $E_{\rm tr}$ equals

$$E_{\rm tr} = \left(\frac{\epsilon_0}{2k_b T_{\rm mwb}}\right) m_p c^2 \sim 3 \times 10^{20} (1+z)^{-1} \text{ eV} , \quad (66)$$

with $\epsilon_0 = m_\pi c^2 (1 + m_\pi/2m_p) = 145$ MeV the threshold energy in the proton rest frame. For energies $E \leq E_{\rm tr}$ particles only interact with photons in the high-energy tail of the Planck distribution, which causes the exponential factor in $l(E)$ (see Berezinskii et al. 1990, chap. 4.4). Associated with this loss length, l_π, one can define a typical energy (see eq. [27]), for $z \ll 1$,

$$E_\pi = ZeB\beta_s^2 l_\pi = 10^{22} Z B_{-6} \beta_s^2 \text{ eV} . \quad (67)$$

If one balances the loss rate (eq. [68]) due to pion production on the microwave background with the acceleration rate at the shock,

$$\frac{dE}{dt} \simeq \frac{ZeB}{c} V_s^2 - \frac{c}{l(E)} E , \quad (68)$$

it follows that losses due to pion production result in a limiting energy for protons of order

$$E_{\max,\pi} \simeq \frac{E_{\rm tr}}{\ln(E_{\rm tr}/E_\pi)} \approx 4.5 \times 10^{19} \text{ eV} . \quad (69)$$

The transition between a situation where the age $t_s \simeq R_s/V_s$ of the shock determines the maximum energy and the case where pion losses determine the maximum energy occurs when $E_{\max} > E_{\max,\pi}$, or

$$R_s \geq \frac{(E_{\rm tr}/ZeB\beta_s)}{\ln(E_{\rm tr}/E_\pi)} \simeq 4.5 B_{-6}^{-1} \beta_{-2}^{-1} \text{ Mpc} . \quad (70)$$

This shows that these losses become marginally important for shocks on scales of ~ 5 Mpc. They are unimportant for smaller shocks.

At these high energies, the discrete character of particle-photon collisions start to play a role. Near threshold for pion production in the proton rest frame, a proton loses an amount of energy per collision given by

$$\frac{\| \Delta E \|}{E} = K_{\pi\gamma} \approx \frac{m_p \epsilon_0}{(m_p + m_\pi)^2} \approx 0.12 . \quad (71)$$

Much above the threshold energy one has $K_{\pi\gamma} \simeq 0.5$. The number of photons encountered by the most energetic particles, which have a residence time $\tau \simeq t_s = R_s/V_s$ at the shock, equals

$$N_{\rm phot} = \frac{R_s}{K_{\pi\gamma} \beta_s l(E)} . \quad (72)$$

This leads to a spread around the maximum energy of order $\Delta \ln E \simeq N_{\rm phot}^{-1/2} K_{\pi\gamma}$, which is

$$\Delta \ln E \simeq \left[\frac{K_{\pi\gamma} R_s}{\beta_s l(E)}\right]^{1/2} = K_{\pi\gamma}^{1/2} \left(\frac{E}{E_{\max,\pi}}\right)^{1/2} . \quad (73)$$

Typically, for particles produced at shocked pancakes on supercluster scales, one has $E_{\max} \simeq E_{\max,\pi}$ and $\Delta \ln E \simeq K_{\pi\gamma}^{1/2} \approx 0.3$.

The luminosity of the pancake is

$$L \sim \frac{M V_{\rm ff}^3}{R} \sim \left(\frac{V_{\rm ff}^5}{G}\right) \sim 2 \times 10^{48} h^{0.5} R_{10}^{0.5} \text{ ergs s}^{-1} . \quad (74)$$

If we further assume that collapsing structures are spaced a distance $r \simeq 50\ h^{-1}$ Mpc apart, then the typical flux at Earth from neighboring pancakes is

$$S = \left(\frac{V_{\rm ff}^5}{4\pi G r^2}\right) \sim 10^{-5} h^{2.5} R_{10}^{0.5} \text{ ergs cm}^{-2} \text{ s}^{-1} . \quad (75)$$

The typical flux in UHECRs inferred from the data of Bird et al. (1993) is (see Shapiro & Silberberg 1983)

$$S_{\rm cr}(E_{\rm eV} > 10^{18.5}) \approx 2 \times 10^{-9} \text{ ergs cm}^{-2} \text{ s}^{-1} . \quad (76)$$

Thus there is enough power in the cosmic rays from pancakes if the efficiency of acceleration of high-energy particles in the shock is of order 10^{-4}. Of course this supposes that structure formation is an ongoing process which is found in many of the standard cosmological models.

To summarize the previous sections, Table 1 lists the possible UHECR production sites, the typical maximum energy E_m that can be produced there, the maximum distance to that source, taking into account redshift due to universal expansion or pion production on the CMWB,

$$D_{\max} = \min\left[\frac{c}{H_0},\ l(E_m)\right] , \quad (77)$$

and the number of sources expected within the sampling volume, given their mean density $\langle n \rangle$. We remark here that only in the case of clusters ($E_m \sim 10^{19}$ eV) $D_{\max} \sim 1$ Gpc is set by electron-positron pair production rather than pion production on the CMWB.

TABLE 1

MAXIMUM ENERGIES

Source	E_m	$D_{\max}$	$\langle n \rangle D_{\max}^3$
AGN	10^{16} eV	~ 3 Gpc	10^5
Jets	10^{20} eV	200 Mpc	~ 10
Cocoons	5×10^{19} eV	500 Mpc	~ 100
Clusters	10^{19} eV	1 Gpc	$\sim 10^4$
Superclusters	5×10^{19} eV	500 Mpc	~ 100

For references, see Table 2.

7. EVENTS ABOVE 100 EeV

There are now seven events above 100 EeV (e.g., Teshima 1993) and three above 200 EeV (Bird et al. 1993; Sommers 1994; Egorov 1994). In Table 1 we summarize our best efforts at obtaining energies of 100 EeV. In some cases, where noted, the numbers have been rather stretched by assuming the existence of large ambient or self-generated fields. However, the overall impression is that there are several possibilities for acceleration of UHECRs up to a few times 10^{19} eV. We continue to assume that these are protonic cosmic rays although at such energies it might even be remotely possible to have the occasional neutrino (Halzen et al. 1995).

The Fly's Eye and Yakutsk events do not apparently emanate from any nearby special object (Ebert & Sommers 1995), although there may be a weak correlation with the direction of the supergalactic plane (Stanev et al. 1995). The lack of any definite confirmed correlation may argue for a more global statistical acceleration process that involves the UHECR component at the highest energies and with the largest gyroradii. This "second-stage" acceleration would have to accelerate particles from $\sim 10^{19}$ to $\sim 10^{21}$ eV. These particles would then be a fundamentally isotropic component, generated by, and interacting with an ensemble of sources such as those listed in Table 1. With this type of model the acceleration would proceed to those energies where the losses due to photodestruction, and not the age of the shock, become the limiting factor. In this section we investigate the viability of such a scheme.

For a second-stage acceleration mechanism to work, two conditions must be met. In the first place, the acceleration rate at the shock must exceed the loss rate due to pion production on the CMWB. This limits the energy of particles for which a second-stage mechanism can operate, given β_s and B, to

$$E \leq E_{\max,\pi} = \frac{E_{\rm tr}}{\ln(E_{\rm tr}/ZeB\beta_s^2 l_\pi)} . \quad (78)$$

Particles with energy exceeding $E_{\max,\pi}$ continue to *lose* energy at the shock at a reduced rate.

Second, the particles must gain more energy at the shock than they lose in transit between encounters with shocks. If the energy gain resulting from the interaction with a shock equals

$\Delta \ln E = f$, and if the average path length between shock encounters is $\bar{l}$, this condition corresponds to $l(E) \geq \bar{l}/f$, or equivalently

$$E \leq \frac{E_{tr}}{\ln(\bar{l}/f l_\pi)} . \tag{79}$$

For shocks with a number density n_s and typical size R_s the path length $\bar{l}$ is equal to the mean free path due to the geometrical cross section of the shocks,

$$\bar{l} \sim l_s \equiv \frac{1}{n_s \pi R_s^2} , \tag{80}$$

provided, of course, that particles propagate scatter free between shocks and are only strongly scattered in the immediate vicinity of the shocks during shock acceleration. If this condition is not satisfied because the scattering mean free path in the IGM is small so that $\lambda_f < l_s$, diffusive propagation between shocks results in a path length of order $\bar{l} \sim l_s^2/\lambda_f > l_s$. Our lack of knowledge of the properties of the intergalactic magnetic field (field strength, coherence length, etc.) precludes any detailed discussion of this point.

In these estimates we have neglected the additional energy loss due to the universal expansion (redshift), $dE/dt = -H(t) E$, with $H(t)$ the Hubble constant. At energies exceeding

$$E \geq \frac{E_{tr}}{\ln(c/H_0 l_\pi)} \sim 5 \times 10^{19} \text{ eV} , \tag{81}$$

the effect of redshift is less important than losses due to pion production.

As long as $\lambda_f \ll R_s$ near the shock, the number of shock crossings is limited by the available time. If the shock expands by an amount ΔR_s, particles increase their energy by an amount of order (see eq. [11])

$$\Delta E \sim ZeB\beta_s \Delta R_s \simeq E_{max} \frac{\Delta R_s}{R_s} , \tag{82}$$

boosting their energy at best a few times above the maximum energy E_{max} for the shock. Thus, at best, we find $\Delta \ln E = f = \mathcal{O}(1)$.

These conditions put strong constraints on the type of shock which could be effective as a second-stage accelerator. Young shocks, or shocks which are continuously driven as in the case of Mach disks in jets, tend to be fast, and therefore $E_{max,\pi}$ tends to be large. But these same shocks are also small, and not very numerous so that $\bar{l} \gg l_\pi$, limiting the energy for which a net acceleration results to values much less than E_{tr}. Old shocks, on the other hand, such as those associated with fossil quasars or (super)clusters, are large, and can have $\bar{l} \sim l_\pi$. But they are also slow, typically $\beta_s \sim 10^{-3}$ to 10^{-2}, resulting in a value of $E_{max,\pi} \leq 10^{18} B_{-6}$ eV. So again, a second-stage acceleration up to energies exceeding 10^{20} eV is not possible.

If one wants to accelerate particles to an energy exceeding 10^{20} eV, the above constraints on second-stage acceleration translate, up to factors of order unity, to the requirements

$$E_\pi \geq E_m \sim 10^{20} \text{ eV} ,$$
$$\bar{l} \leq l_\pi \sim 10 \text{ Mpc} , \tag{83}$$

for the population of shocks responsible for the acceleration. For protons ($Z = 1$) these conditions can be written as

$$\beta_s^2 B_{-6} \geq 10^{-2} ;$$
$$\Upsilon \equiv n_s \left(\frac{c}{H_0}\right)^3 \geq 10^7 h^{-3} R_{10}^{-2} , \tag{84}$$

TABLE 2

SOURCE PARAMETERS

Object	$\beta_s^2 B_{-6}$	R_{10}	Υ
Hot spots[a]	~ 100	10^{-4}	10^5
Cocoons[a]	$\leq 10^{-2}$	10^{-2}	10^5
Quasar bubbles[b]	$\sim 10^{-6}$	~ 1	5×10^7
Clusters[c]	$\leq 10^{-6}$	~ 0.5	3×10^5
Superclusters	$\leq 10^{-4}$	~ 3	10^4

[a] Typical density of powerful radio galaxies from Burns, White, & Hough 1981.

[b] Typical quasar density from Chokshi & Turner 1992, assuming a median black hole mass of $10^8 M_\odot$.

[c] Canonical numbers from Bahcall & West 1992.

where $R_{10} = R_s/(10 \text{ Mpc})$. The quantity Υ corresponds to the number of shocks per cubic Hubble radius. As Table 2 shows, no objects are capable of satisfying both these requirements simultaneously.

We are therefore forced to conclude that no known object can be responsible for second-stage acceleration of particles to energies exceeding 10^{20} eV. Therefore, particles above 100 EeV must have been accelerated in "one shot" at their parent objects and should be identifiable using giant arrays. As can be seen from the calculations in the previous sections, and from Table 1, the objects responsible for these particles must be rather extreme in the range of possibilities.

8. SUMMARY AND PREDICTIONS

We have given a comprehensive study of all possible conventional sources that can explain the protonic component of UHECRs above 3 EeV and extending up to 300 EeV. Our results are summarized in Tables 1 and 2. We have argued that cosmic rays in the energy range $E > 10^{16}$ eV cannot be produced in, or escape from escape from AGNs. Using a metallicity-composition argument which applies for the spectrum above and below 30 EeV we have excluded any origin for UHECRs above $10^{18.5}$ eV in metal-rich environments such as rich clusters, the inner regions of galaxies, and extended halos. Therefore, we have sought the origin of these particles in large-scale extended structures associated with the relatively unenriched component of the intergalactic medium.

Large (Mpc-)scale shocks from explosions and winds, still active within ~ 1 Gpc, can generate 10^{20} eV protons if they self-generate microgauss fields. Both hot spots and cocoons of powerful (F-R Type II) radio sources are plausible sources, with the principal uncertainty in both cases being the magnetic field strength that is adopted or inferred. Cocoons associated with extragalactic jets are a plausible source for UHECRs up to $10^{19.5}$ eV.

Cosmic shocks, formed as structure develops in the gravitational collapse of primordial perturbations as seen in the standard cosmological models such as in pancake-like structures, and in the collisions of hierarchical merging subunits, are also possible sites for the UHECR acceleration up to $10^{19.5}$ eV, provided there is a primordial field of $\gtrsim 10^{-9}$ G or if microgauss fields can be self-generated in shocks.

We have analyzed in detail the possibility of a second-stage acceleration of the highest energy cosmic rays by an ensemble of shocks but as is summarized in Table 2 we have been unable to identify a sufficiently powerful class of shocks with a sufficiently large volume-filling factor that could achieve a diffusive acceleration in a large-scale shock network pervading the IGM. Thus we conclude that particles above 100 EeV should come from their parent sources and their origin in space should be identifiable by the currently planned giant arrays.

Large-scale shocks that generate relativistic particles should be relatively common (Table 2) throughout the local universe. Although interesting possibilities for cosmic-ray acceleration from topological defects have been proposed by Sigl et al. (1994), there may not be sufficient flux from such sources (Gill & Kibble 1994). Conventional shock acceleration in admittedly extreme but plausible situations can account for the current observations. These considerations indicate that the IGM has a rich structure of shocks if the old remnant shocks from the galaxy formation and QSO epochs are included. Serious attempts should be made to observe them both in colliding substructures where the shock is already indicated by the X-ray observations and filamentary and pancake structures. In addition, large-scale remnant shocks from active galaxies, quasars, and starbursts and their remnants should be sought at very faint surface brightness. Many important parameters in the current analysis would be more significantly constrained by such observations including the magnetic field strength and particle acceleration efficiency in such structures.

We have not considered the UHECR propagation through the IGM and the observed isotropy of arrival directions here in any detail (Achterberg et al. 1995). The UHECR energy flux (eq. [73]) corresponds to an associated power for a mean source distance D of order

$$L_{\rm cr}(>10^{18.5}\ {\rm eV}) = 4\pi D^2 S_{\rm cr}(>10^{18.5}\ {\rm eV}) \approx 2\times 10^{41} D_{\rm Mpc}^2\ {\rm ergs\ s^{-1}}\,. \quad (85)$$

Thus, say, several sources at ~ 10 Mpc of luminosity $L \sim 10^{45}$ ergs s^{-1} are required with an efficiency of production of $\sim 10^{-2}$. After this paper was completed we have learned of two studies of the interesting possibility that gamma-ray bursts may be producers of the highest energy cosmic rays (Milgrom & Usov 1995; Waxman 1995; Vietri 1995). As clearly discussed in detail by Waxman, obtaining the flux is the most severe problem since it is apparent that the efficiency of UHECR production in the fireball must be rather high ($\sim$0.1–1), which, although not out of the question, is several orders of magnitude more efficient than the more conventional shock wave acceleration mechanisms analyzed here.

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