ACCELERATION OF ULTRA HIGH ENERGY COSMIC RAYS: COSMIC ZEVATRONS?

T.W. JONES
Department of Astronomy
University of Minnesota, Minneapolis, MN, 55455 USA

Abstract.
In this lecture I outline some of the underlying physics issues associated with accelerators plausibly capable of explaining the UHECRs up to ZeV energies. I concentrate on the mechanisms and their constraints, but provide a brief background on observations and the constraints they supply, as well.

1. Introduction
Strong observational evidence now exists for the existence of cosmic rays (CRs) at energies approaching 1 ZeV (=10^{21} eV). Their origins remain a mystery and a serious theoretical challenge. In this lecture I will outline some of the key issues that constrain theories for production of these “Ultra High Energy Cosmic Rays” (UHECRs). I will focus on so-called “bottom-up” scenarios, which accelerate charged particles, usually protons, from low energies. These hypothetical accelerators have been dubbed “Zevatrons”, since they must achieve ZeV energies. Other lecturers at this School have addressed the alternative “top-down” scenarios, which explain UHECRs as decay products of fossil Grand Unification defects, or by “new” physics (e.g., Stecker & Sánchez). The new physics approaches also include explanations that reduce constraints on long range UHECR propagation by allowing violation of Lorentz invariance, so those approaches can combine both scenarios. In practice, both bottom-up and top-down scenarios are generally forced to fairly extreme model assumptions in order to account for current measurements. So, which ever approach turns out to be most relevant, we will learn some very interesting physics or astrophysics. The coming generations of cosmic ray detectors should tell us which direction to go within the next few years.
In addition to the other lectures from this School readers interested in more details or alternative perspectives may want to read some of the excellent reviews of this topic that have appeared recently in the literature (e.g., Olinto 2000; Watson, 2000; Sigl, 2002).

I begin my discussion with a short outline of some background that helps to define the underlying issues. My main goal is to confront the physics that must be included in cosmic Zevatrons, so §3 aims to find the underlying constraints that can be used to limit the field of models. Since many models accelerate particles in shock waves, I will follow with a short discussion of the physics of diffusive shock acceleration, before turning to apply the appropriate physical constraints to some candidate phenomena in §5.

2. Some Background

The cosmic ray energy spectrum has been measured over almost 13 decades in energy, roughly between \( \sim 0.1 \text{GeV} \) and \( 10^{12} \text{ GeV} (= 1 \text{ ZeV}) \). At lower energies the solar wind traps extra-solar-system charged particles and carries them away. At higher energies the extrapolated fluxes are extremely small at best \((< 10 \text{ particles km}^{-2} \text{ century}^{-1})\), so not yet measured or adequately constrained. To a first approximation the measured energy spectrum is a broken powerlaw with a slope close to -2.7 below a “knee” near \( 10^6 \text{GeV} \), then steepening by about 0.5 in the index above the knee. The spectrum appears to flatten slightly again above \( \sim 10^9 \text{GeV} (= 1 \text{ EeV}) \). This feature is sometimes called the “ankle” of the CR spectrum. The knee and ankle features are commonly viewed as corresponding to transitions either in the nature of the CR accelerator or in the propagation of the CRs. As outlined below there is currently a controversy about the spectrum above \( 100 \text{ EeV} \) (Abu-Zayyad, et al. 2002; Takeda, et al. 2002), but several groups have independently detected events with energies in this range, so there is little doubt about their existence. The highest energy event reported so far is \( 0.3 \text{ ZeV} \) (Bird, et al. 1994). That amounts to an astounding 50 Joules of kinetic energy in an individual elementary particle. Not only is it problematic to explain the generation of CRs at these high energies, their propagation is severely limited, especially by interactions with the cosmic microwave background (CMB). The difficulties and implications are obviously very important, so there are major efforts underway to improve on current measurements.

The composition of the cosmic rays near or above the knee is not well established. As discussed by Stecker in his contribution at this School (Stecker 2002) the evidence mostly supports a transition from a mix of light and heavy nuclei near and just above the knee to mostly light nuclei at the highest energies. Most, but not all discussions of UHECRs take them to be
protons for that reason. Additionally, collisions between ultrahigh energy heavy ions and photons disrupt heavy ions, so that they cannot propagate over long range. (See Stecker 2002 for a discussion of these properties.)

Another important property of the UHECRs that constrains their origins or possibly their propagation is the distribution of detected events on the sky. Below the knee CR detections are nearly isotropic (e.g., Erlykin, et al. 2002), but that is a reflection of that fact that low energy CR propagation is diffusive in the turbulent galactic magnetic field. Sources of these CRs are obscured because of that. On the other hand, since the gyroradius of a particle with charge \( Z \) and energy \( E_{\text{eV}} = E / \text{EeV} \) is \( r_g \approx (E_{\text{eV}} / Z B_{\mu G} \text{ kpc}) \), and the characteristic galactic magnetic field is a few \( \mu \text{G} \), confinement of CRs within the galaxy is limited to energies below \( \sim 1E_{\text{eV}} / Z \). Diffusive propagation within the galaxy should break down well below that energy. Indeed, around the ankle some enhancement in the CR flux is reported towards the galactic center (Hayashida, et al. 1999), so we may have seen direct indications of their galactic origins. At higher energies, however, the data are consistent with isotropic detections (Takeda, et al. 1999; Uchihori, et al. 2000). Unless the UHECRs are highly charged \( (Z \gg 1) \) this almost assures their extragalactic origins.

2.1. SOME PROPAGATION ISSUES

Before discussing the physics of various acceleration scenarios it is helpful to establish a clear picture of what general limits can be placed on the locations of possible extragalactic sources of the UHECRs. The most frequently discussed constraint on the origins of the observed particles comes from the large energy losses by hadrons through inelastic collisions with photons. This is also an important limit to acceleration models, as we shall see. Collisions between protons and photons can generate leptonic pairs if total energy in the center of momentum frame exceeds the proton mass by more than the pair mass. The proton energy threshold, \( E_{\pm} \) for this photopair production when the incident photon energy is \( \epsilon_\gamma \) is approximately \( E_{\pm} \approx 10^3 / E_\gamma (\text{eV}) \text{ GeV} \). For the CMB \( \epsilon_\gamma \sim 10^{-3} \text{eV} \), so the photopair production threshold energy is about 0.01 \( \text{EeV} \).

The losses are much more serious, however, if the collision energy is enough to excite the \( \Delta^+ \) resonance near 1.2 GeV in the center of momentum frame, since that decays to produce pions, typically carrying away of order 10% of the total energy. This threshold is roughly \( E_\pi \sim 5 \times 10^7 / \epsilon_\gamma (\text{eV}) \) GeV. For CMB photons the effective threshold is just below 100 \( \text{EeV} \). The cross section rises sharply, so that protons near the threshold lose most of their energy in the CMB in less than a few \( \times 10^8 \) years. Thus their effective propagation length is less than about 100 Mpc. The significance of this was
pointed out in the 1960s independently by Greisen (1966) and by Zatsepin & Kuz’min (1966). If UHECRs are of cosmological origins the result should be a relative absence of particles above \(\sim 100\) EeV. If the source spectrum is sufficiently flat above this “GZK” feature, one might also expect a flux pile up around 100 EeV as well. The experimental situation is unclear at the moment, since groups using complementary techniques for analyzing extensive air showers have obtained conflicting spectral properties above 100 EeV. The HiRes experiment, using atmospheric fluorescence measurements seems to see the GZK feature (Abu-Zayyad, et al. 2002). The AGASA muon detector reports a spectrum that actually hardens towards the highest energies (Takeda, et al. 2002). I will consider this an open question in the following discussions. The Auger Observatory, which will apply both techniques simultaneously and will have a substantially enhanced effective target area, should resolve this question in the next several years, in any case.

A related propagation issue is whether the UHECR events point back close to their sources; that is, whether intergalactic magnetic fields have any appreciable influence on UHECR trajectories. This obviously determines the degree to which the observed distribution mimics the source distribution. Magnetic fields in extragalactic space are not well known. In some rich clusters there is evidence for fields of at least several tenths of a micro Gauss, but elsewhere they should be much weaker (e.g., Kronberg 2001). The gyroradius can be written conveniently as \(r_g \approx E_{\text{ZeV}} / (ZB_{\text{nG}})\) Gpc, where B is now expressed in nano Gauss. Except in clusters or possibly the vicinity of galaxies the expected deflections should be small over distances that hadronic CRs can propagate without severe energy losses. Therefore most UHECR explanations anticipate that detected events point back to the original sources, which must then be relatively isotropic. Some analyses have suggested enhancements in the UHECR distribution towards “interesting” regions of the sky, particularly the super galactic plane, but at present the evidence does not seem to support any such biases (e.g., Takeda, et al. 1999; Uchihori, et al. 2000).

The possibility that particle trajectories are significantly deflected is still an open question, however (e.g., Medina-Tanco, et al. 1997; Lemoine, et al. 1997; Biermann, et al. 2001; Medina-Tanco & Enßlin 2001). Biermann, et al. 2001, for example, have argued that a wind from our galaxy could carry a strong enough magnetic field to sufficient distance that UHECR trajectories would be strongly modified. They have shown that such a field might redistribute the 13 highest energy events detected at that time in a way consistent with their common origin being within the Virgo cluster. (See the lecture in this volume by Biermann for more details.) On the other hand, the existence of this field structure is not established by other means,
so it remains hypothetical. We may hope that greatly improved UHECR statistics in the next several years will either define unique event patterns or exclude them from consideration.

Incidentally, strong deflections of particle trajectories also add to the path lengths, so increase the influence of the photopion production losses above the GZK feature unless the sources are particularly close by. (Virgo is only about 18 Mpc away, so this is not much of an issue in that case.) Since the path lengths become energy dependent, it would also spread out the arrival times of simultaneously emitted particles with the highest energy particles arriving first.

There is at least one other basic constraint we can establish for the UHECR sources from the limited information already established. The measured flux combined with the energy loss rates establishes an estimate for the energetics of the associated phenomenology. The flux at 100 EeV translates into a local energy density of about $3 \times 10^{-22}$ J m$^{-3}$. Taking the loss time as $3 \times 10^8$ yr and assuming a steady state, the mean luminosity of UHECRs per unit volume must be about $10^{-37}$ W m$^{-3}$ or about $3 \times 10^{44}$ erg Mpc$^{-3}$ yr$^{-1}$. That is comparable to the estimated mean volume luminosity of Gamma Ray Bursts, a fraction of one moderate Active Galactic Nucleus (AGN) inside 100 Mpc, or the power dissipated by the accretion shock of a cluster like Coma or Perseus, as examples of phenomena that have been called upon to explain UHECRs.

In summary, it seems likely that UHECRs are extragalactic in origin. If the GZK feature is not evident then the dominant sources should be within about 100 Mpc. If the feature is seen, and especially if there is a pile up around 100 EeV, then a more extended distribution of sources becomes more likely. There are a variety of phenomena that may be energetic enough to produce UHECRs. It is unclear yet if the distribution of events on the sky represents the distribution of sources or has been significantly modified by extragalactic magnetic fields. With these factors in mind let now look at some of the basic physics constraints that must be satisfied by successful accelerators of UHECRs.

3. How to Make a Cosmic Zevatron

Accelerating a charged particle to ZeV energies is not easy! We can express how difficult it is in a way that is almost independent of the detailed physics of the acceleration process. I will illustrate this point with several lines of reasoning that lead to a common relationship first emphasized by Hillas (1984).

The simplest constraint would be that the accelerated particles must be contained inside the accelerator while they are being accelerated. So,
we must demand that the particle gyro radius be less than the size of the accelerator. That is

$$r_g = \frac{pc}{ZeB} = \frac{E}{ZeB} < R,$$  \hspace{1cm} (1)

where \(E\) is the particle energy, or

$$E < ZeBR = 0.9ZB_G R_{pc} Z \text{eV}$$  \hspace{1cm} (2)

where \(R\) is a characteristic size for the accelerator.

Magnetic fields cannot accelerate charged particles; only electric fields can do that. It is not always easy to express the physics of the accelerator in terms of an applied electric field, but if it is, then obviously the maximum energy that can be achieved is given by \(E < ZeER\), where \(E\) is the electric field. For a unipolar inductor, for example, \(E \sim v_{acc} B/c = \beta_{acc} B\), where \(v_{acc}\) is a characteristic large scale speed (e.g., \(\Omega \times R\)). This argument leads to the expression

$$E < \beta_{acc} ZeBR,$$  \hspace{1cm} (3)

which differs from the relation (2) only by the velocity factor, \(\beta_{acc}\).

A related argument is that the time necessary to accelerate a particle must be less than the lifetime of the accelerator. Let us apply this to the diffusive shock acceleration process to be outlined in §4. The mean time to accelerate a particle to energy \(E\) is given by equation 9, so we have

$$t_{acc} \sim \frac{C_d \kappa}{v_{shock}^2} < t_{dynamic} \sim \frac{R}{v_{shock}},$$  \hspace{1cm} (4)

where \(\kappa\) is a characteristic spatial diffusion coefficient for the highest energy particles to be accelerated, and \(C_d \sim 10\) is a numerical factor. Under most circumstances the applicable diffusion coefficient is \(\kappa \sim (1/3)\lambda c\), where \(\lambda\) is the scattering length of the particles. If we require that the scattering length exceed the particle gyroradius we once again obtain relation (3) above within a numerical factor of order unity.

I will simply call relation (3) the “Hillas constraint”. As did Hillas, we can quickly filter phenomena from our discussion of viable UHECR candidate sources with a plot of this relation. The line in Figure 1 shows the relationship between source size and magnetic field that would be given by an equality in relation (3) for a particle of energy \(100\text{ EeV}/(Z\beta_{acc})\). For comparison I have marked rough coordinates for some energetic phenomena that are thought likely to produce cosmic rays in some form. By the arguments given, only points above the line are viable candidates to produce \(100\text{ EeV}\) CRs. This excludes supernova remnants or winds from starburst galaxies, for example. They simply do not have the combined size and magnetic field to produce UHECRs.
There are further basic filters that must be applied to phenomena passing this first test, starting with limitations from energy losses within the accelerator. In particular the energy loss times must obviously exceed the maximum of the acceleration time and the escape time from the accelerator. The two energy loss processes that most commonly limit acceleration to the highest energies are synchrotron emission and photon scattering. Inverse Compton scattering (elastic photon scattering) looks much like synchrotron emission, so is usually included in that context. Synchrotron/Compton energy losses scale as $E^2$, so become smoothly more important as the energy increases. The energy loss rate from inelastic collisions rises sharply in two stages due first to photopair production and second to photopion produc-
tion, then remains roughly constant, with the thresholds scaling with $\epsilon^{-1}$ as described in §2 (see also, for example, Stecker 2002, Berezinsky, et al. 2001).

The energy loss lifetime of a particle from synchrotron/Compton emission is simply

$$t_{s,c} = \frac{3}{2} \frac{Z^4 e^4 (B^2 + B_{eq}^2)}{m^4 c^7} \frac{1}{E^7}$$

where $m$ is the particle mass and $B_{eq} = \sqrt{6 \pi \mu_{rad}}$ is sometimes called the equivalent isotropic magnetic field for the photon field in terms of Compton losses. For a black body radiation field $B_{eq} = 0.38 T^2 \mu G$.

Inelastic scattering losses are usually dominated by photopion production, which turns on roughly above $E_\pi \sim 100$ EeV/T for a black body radiation field. We can approximately model the energy loss lifetime above this threshold as

$$t_\pi \sim 3 \times 10^{16} \frac{1}{T^3} \sec.$$  

The simplest constraint, with the escape time set to the free flight time, $R/c$, is

$$\min(t_\pi, t_{s,c}) > \max(t_{acc}, R/c).$$

These limitations can become effective under a variety of circumstances, but especially when the particle accelerator is associated with an intense magnetic field or radiation field (i.e., “hot” in the black body sense) or if the escape time or acceleration time is long. The former condition applies, for example, to AGNs or to young pulsar magnetospheres. In a $10^7$ K radiation field around a young neutron star TeV protons can propagate only a few km, so cannot escape the magnetosphere or be accelerated to still higher energies (Venkatesan, et al. 1997). Similarly the radiation field near an accreting massive black hole in an AGN will limit the escape energy of protons to less than a few hundred TeV. In either case, accelerator regions further out from the central massive object remain viable by this condition.

Models of Gamma Ray Bursts (GRBs) typically call for very strong magnetic fields in equipartition with the radiation field (e.g., Mészáros & Rees 1994). Then synchrotron losses are particularly important and can provide a strong limitation to the energies of escaping particles, which may still come close enough to ZeV to make them interesting. Adiabatic losses would also be severe if the acceleration took place only at very early times in the expansion of the fireball, although conditions may be suitable long enough to avoid this issue.

At the other extreme of source environments, photopion energy losses probably eliminate cosmic structure shocks as viable candidates for UHE-CRs, at least above the GZK feature (Norman, et al. 1995; Kang, et al. 1996; Kang, et al. 1997). The difficulty here is the long time required for particles
to be accelerated to these energies. These shocks have characteristic inflow speeds of only \( \lesssim 0.01c \), and the magnetic fields are probably \( \lesssim 1 \mu G \). Those lead to acceleration timescales of at least a few times \( 10^8 \) years near 100 EeV, and that will be recalled as the limiting propagation time through the CMB associated with the GZK feature.

These various constraints have significantly narrowed the field of candidate phenomena, but there are still several interesting possibilities. On the other hand we have so far only established constraints and said nothing about the physics of the acceleration process itself. In the next section I will remedy this, focusing especially on acceleration at shocks, since that mechanism is the one most commonly invoked to produce high energy charged particles.

4. Particle Acceleration Physics

4.1. GENERAL COMMENTS

Cosmic Rays are usually discussed as a distinct population of particles in the universe, although that is largely a consequence of the distinct ways they reveal themselves to us. This approach to the physics is certainly appropriate if UHECRs, for example, are byproducts of the decay of relic massive particles or accelerated by large scale electric fields, in pulsars, for instance. However, a somewhat more realistic way to view these particles under most circumstances is as a high energy component of the ionized matter in the universe. Galactic CRs, for example, have roughly the same elemental composition as the interstellar medium at those energies where this property is well measured. They are very likely to have come from the interstellar medium.

Cosmic plasmas are largely “collisionless” in the sense that binary Coulomb collisions play a minor role in dissipation of energy. Instead, dissipation comes from “collective” interactions that depend on fluctuations in charge and currents (i.e., plasma waves or turbulence). These interactions tend to be selective in their efficiency, so that cosmic plasmas are probably rarely fully thermalized. The problem is that they never have the time and they are not really closed systems in the thermodynamic sense. This especially applies to the most energetic charged particles, a point that was also central to our discussions in the previous section. In that sense cosmic rays can be seen as those particles that have so-far escaped being thermalized, but have preferentially extracted energy from a system out of thermodynamic equilibrium. The prototype acceleration model for this was that devised by Fermi (1949). There are several modern variants of his idea; they all involve stochastic processes in which charged particles with effective mean free paths too long to allow them to become thermalized extract free energy
via interaction with fluctuations in a dynamic system. The currently most commonly discussed acceleration process, diffusive shock acceleration, is of this type, and I will devote some space to its details momentarily. I note that not all collisionless plasma models depend on Fermi processes, however. In some other settings, such as magnetic reconnection, acceleration is often predicted as a result of locally strong induced electric fields. However, The microphysics of this once again involves the details of a breakdown in thermodynamic or bulk magnetohydrodynamic behaviors (e.g., Lesch & Birk 1998).

4.2. DIFFUSIVE PARTICLE ACCELERATION AT SHOCKS

Starting from the premise above, we can understand so-called diffusive shock acceleration (“DSA”) as a part of the physics of collisionless plasma shocks. Cosmic shocks would mostly not exist if they had to depend on binary interactions; they form, instead due to very complex, nonlinear phenomena involving plasma turbulence (e.g., Quest 1988). The end result is a transition that in some respects can be described by classical fluid shocks, but in others may not be. Like classical shocks a region develops where most of the particles entering from upstream interact strongly and have much of their directed energy “dissipated”; that is they are “thermalized”. In this context we will identify CRs as those particles that are not thermalized because they interact too weakly with the plasma turbulence to become trapped inside the primary dissipative layer. The dissipative layer is usually called the shock, although it is only part of the shock structure, so I will use the more appropriate “subshock” label. The thickness of the subshock is not easily characterized, and it depends on the orientation of the local mean magnetic field. Often it is described in terms of the characteristic gyroradius of thermalized particles.

If we can ignore the mass, energy and momentum carried by the CRs then the usual hydrodynamic or MHD shock jump conditions apply across the subshock. CRs then can be treated as test particles interacting with this flow. The central idea behind DSA is that these CRs still are sufficiently strongly scattered by local plasma turbulence that their motions become randomized (but not thermalized) close to the subshock. The resulting behavior is fairly easily described if the speed of the CRs is large compared to the bulk motions relative to the shock and if we can approximate the distribution function, \( f(p) \), as almost isotropic in both the upstream and downstream frames of reference. In that case, for example, the flux of CRs crossing the subshock from upstream in a momentum range \( dp \) is approximately \( \pi v^2 dp f(p) \), where \( v \) is the particle speed. On the other hand, the CRs downstream are advected away at the bulk flow speed of the post-
shock plasma, which we can call $u_2$. The flux there is $4\pi u_2 p^2 dp f(p)$. In a steady state the ratio of these two fluxes determines the probability that a particle entering the shock from upstream will escape downstream at the same momentum, $P_e$. Since the steady downstream distribution function must be uniform and also cannot change across the subshock (the CRs ignore the subshock, remember), this gives $P_e \approx 4u_2/v << 1$. Thus, under these assumptions most of the CRs are temporarily trapped between the upstream and the downstream flows. Also, by construction they are not stopped within the subshock itself, so that they are akin to ping pong balls trapped between converging paddles, here represented by the large scale plasma turbulence. It is simple to show that on average each time they cycle from upstream to downstream and back again they gain momentum given by

$$\frac{\delta p}{p} \approx \frac{4}{3} \frac{u_1 - u_2}{v} = \alpha << 1,$$

where $u_1$ represents the upstream bulk plasma speed into the shock. Once the particles are relativistic the fractional momentum gain, $\alpha$ is a constant, as is the escape probability, $P_e$.

Under such circumstances Fermi showed that the steady state momentum distribution is a powerlaw (since there is no associated momentum scale) given by $f(p) \propto p^{-q}$, with $q = 3 + \frac{P_e}{\alpha}$. Fermi was not thinking of shocks, but in the late 1970s several people independently recognized how this Fermi-like process would work in shocks (Axford et al. 1977; Krymskii 1977; Bell 1978; Blandford & Ostriker 1978). In that case the powerlaw index of the momentum distribution becomes $q = \frac{3r_1}{r_1 - r_2} = \frac{3\rho_1}{\rho_1 - \rho_2}$, where $r = u_1/u_2 = \rho_2/\rho_1$, is the compression through the shock. For a strong, nonrelativistic adiabatic shock $r \rightarrow 4$, so $q \rightarrow 4^+$. For even moderately strong shocks the powerlaw slope is close to 4, so the solution appears pretty robust. Notice that this result depends only on the jump conditions across the shock, and it was necessary only to suppose that somehow plasma turbulence would isotropize the CRs on each side of the shock. Plasma shocks were long known to be highly turbulent downstream. The other key insight to the introduction of DSA was that unless the CRs returning to the upstream side of a shock where isotropized with respect to the upstream flow they would amplify local Alfvénic turbulence there by a streaming instability. Those waves should then, in turn scatter the subsequent CRs trying to escape in this way. This physics, even though it leads to a simple outcome, actually shows that DSA is a nonlinear aspect of shock formation, so was an early clue that the physics is actually more complex.

An important detail about DSA that we encountered in §3 was the finite amount of time required to accelerate particles up to a specified energy (or momentum). This is now easy to estimate from the fractional momentum
gain per shock encounter expressed in equation 8 along with an estimate of the time interval between CR shock crossings. The CRs propagate approximately a diffusion length, \( x_d = \kappa / u \), on each side of the shock before their return. Since the CRs move at a speed \( v \), the average time between a pair of shock crossings is \( 2(x_d(1) + x_d(2)) / v \), where the subscripts refer to the upstream and downstream values for \( x_d \). From this the mean acceleration time is (e.g., Lagage & Cesarsky 1983)

\[
    t_{acc} = \frac{3}{u_1 - u_2} \int_{p_{min}}^{p_{max}} \left[ \frac{\kappa_1}{u_1} + \frac{\kappa_2}{u_2} \right] \frac{dp}{p}
\]

(9)

This expression has been verified by a variety of numerical approaches. In particular it provides a reasonable estimate of the upper cutoff to the CR spectrum one expects if the particles are injected locally at the shock.

The most obvious nonlinear complication in DSA physics comes from a calculation of the pressure associated with the CRs. That is

\[
    P_{cr} = \frac{4}{3} \int_{p_{min}}^{p_{max}} v p^4 f(p) d\ln p.
\]

(10)

It is clear that for the strong shock limit \( (q = 4) \) \( P_{cr} \) diverges logarithmically with \( p_{max} \). Now, that momentum will in practice always be limited by some constraint, such as those we have discussed. Still, it is apparent that the pressure in CRs can potentially compete with the thermalized plasma or even dominate it. This is quite an important realization, especially since the CRs diffuse upstream from the subshock a characteristic length \( x_d \) as part of the acceleration process. That means there will be a pressure gradient \( \sim P_{cr}/x_d \) influencing the flow upstream, so that the flow into the subshock will in this “shock precursor” be fractionally reduced by an amount \( \delta u / u_1 \sim P_{cr} / (\rho_1 u_1^2) \). The maximum postshock thermalized pressure is \( \frac{3}{4} \rho_1 u_1^2 \), so if the CR and thermal pressures expected at the subshock are comparable, the shock must be significantly modified from the classical hydrodynamic shock.

The CR feedback generally tends to enhance the dynamical role of the CRs, since the deceleration through the precursor weakens the subshock and thus reduces the postshock thermal pressure. At the same time it turns out that the total compression through the shock, including the precursor and subshock generally exceeds that for a purely hydrodynamical shock. Then the expected value of \( q \) is reduced for higher energy CRs, which tend to have longer scattering lengths, so see a larger compression. That adds to the CR pressure as well, since the relative number of the most energetic particles is greater. At the same time, if the highest energy CRs escape upstream into a region without sufficient scattering to return them they remove energy (and total pressure) from the shock. This likely eventuality further enhances the
total compression through the shock, much like what happens in a radiative shock. On the other hand, some models explaining how particles manage to be “injected” into the CR population by escaping upstream through the shock for the first time would predict a reduced injection efficiency once the subshock is weakened. That, in turn, would reduce the expected $P_{cr}$, so provide feedback of the opposite sign. Malkov, et al. (2000) have shown a bifurcation in the steady state analytic solutions when these various features are included and argue that a critical, self-organization may develop that determines the weights of the various nonlinear effects. Figure 2, from Kang, et al. 2002, illustrates some these features in the evolution of a simulated CR shock.

There are still other nonlinear features of the DSA process, most of which are less well understood than what I just described. I refer interested readers to some of the recent reviews for further information (e.g., Malkov & Drury 2001; Jones 2001).

Especially in some phenomena considered for acceleration of UHECRs the flows into the shocks may be relativistic or even ultrarelativistic. This changes important details about DSA that make its results harder to calculate. First, the assumption that we could consider the distribution function, $f(p)$, to be isotropic in both the upstream and downstream frames obviously cannot apply, since the bulk speeds are no longer a small fraction of the CR speed. Second, an unstated detail in my previous description was that the orientation of the large scale magnetic field relative to the shock normal could be ignored to a first approximation. However, especially for relativistic shocks an oblique field may intersect the shock face “superluminally”; that is, the intersection moves along the face at superluminal speed. Particles cannot then re cross the shock from downstream by following an individual field line. In that case only motions of particles across field lines are relevant to DSA and that is more difficult to model. Despite these difficulties, the test particle form of the process is reasonably well understood now, and is nicely discussed in recent literature (e.g., Achterberg, et al. 2001). Here I mention only a couple of important comparisons with nonrelativistic DSA properties. On the face of it one might expect DSA to be somehow more efficient in a relativistic shock, since the fractional energy gain per crossing is large, $\propto \Gamma^2$, where $\Gamma$ is the Lorentz factor of the shock. On the other hand, the compression ratio of a highly relativistic shock is less than that for a nonrelativistic shock, asymptoting to 3 rather than 4. This enhances the relative escape probability of the CRs, so leads to a steeper momentum distribution. For asymptotically strong relativistic shocks the expected value of $q \approx 4.2$ rather than $q \approx 4$. Thus, compared to nonrelativistic shocks we expect reduced nonlinear feedback of from CRs (still defined as those particles that are not thermalized).
Figure 2. Evolution of a simulated CR modified shock with initial Mach number 30. The upper four panels show evolution of dynamical variables, including gas density, pressure and flow velocity, along with the CR pressure. The initial conditions were a simple gasdynamic shock, visible as a step function in the plots. The bottom two panels show the evolution of the proton momentum distribution, shown here as $g(p) = p^4 f(p)$. The spatial diffusion coefficient was Bohm diffusion and the times are given in units of the diffusion time for $p = mc$. 
Before returning to discuss astrophysical models for UHECR accelerators it is helpful to mention another proposed acceleration mechanism that is similar to DSA, but which would operate across thin shear layers rather than across shocks. It was pointed out by Jokipii, et al. (1989) that diffusion of CRs across a shear layer would produce similar results to a shock provided the CR mean free paths exceed the thickness of the shear layer. Once again the CRs find themselves caught in an apparently convergent flow. In nonrelativistic flows the process is relatively inefficient, but Ostrowkski (2000) has argued that the large boosts that would accompany this process in the boundary of a relativistic jet might make it astrophysically interesting.

5. Comments on Some Candidate Zevatron Models

Our discussions in §3 considerably limited the range of candidate UHECR accelerators. While there are currently a variety of UHECR accelerator scenarios on the market with serious supporters, I will focus here on only three that are illustrative. The most viable candidate in my view may be high powered jets from AGNs interacting with the intergalactic medium (IGM). Gamma Ray Bursts may be excluded on the basis of their redshift distributions, as I will mention, yet they may be physically just capable of producing such particles, so I will include them. Since the possibility that UHECRs are heavy ions is not yet excluded, I will also include one example of such models; namely, acceleration of Fe ions in young pulsar winds.

5.1. GIANT RADIO JETS

AGNs are thought to derive their tremendous energy supplies from accretion onto supermassive black holes. We excluded the AGN environment itself from the UHECR discussion on account of the severe losses expected from photopion production there. However, a significant subpopulation of AGNs expel large amounts of energy in directed outflows, typically identified through associated radio emissions. These so-called radio jets are probably formed as relativistic beams near the black hole and then drill their way to large distances. In some cases the interaction regions extend as much as a Mpc from the galaxy itself. The energy of the jet is dissipated in a bow shock inside the IGM and in shocks within the jet plasma itself. The bow shock is probably not fast enough to produce UHECRs. In the cartoon model of radio jets there is a strong termination shock at the end of the jet. The characteristic size of such a shock would be $R \sim 10\text{kpc}$; the magnetic field (from radio and X-ray observations) might be $B \sim 10 - 100\mu\text{G}$, and $v_{\text{shock}} > 0.1c$. With these numbers the Hillas constraint for protons would be $E_{\text{max}} \sim 1\text{ZeV}$ (e.g., Rachen & Biermann 1993). The acceleration
timescale $\sim 10^7$ yrs, which is comfortably within the expected lifetime of
the jets. The constraints from synchrotron and photopion losses give similar
limits. Simulations show that the cartoon picture of radio jets is too simple
(e.g., Tregillis et al. 2001), since the termination shock is unsteady and
much of the dissipation actually involves a sequence of somewhat weaker
shocks. Nonetheless, the shocks in the jet are large and very fast, so the
radio jet scenario remains viable in my view. In addition, Ostrowski has
argued that if the relativistic jet has a thin boundary layer particles may
be accelerated quickly by the diffusive shear acceleration process mentioned
in §4. I am not convinced that such thin layers will exist, but perhaps this
could contribute. Other authors have suggested that magnetic reconnection
associated with the jet flows could also play a role (Schopper, et al. 2001).

One important detail here concerns the composition of radio jets. It is
unclear if there is a substantial proton component in these jets. Energy
requirements are lower if the jets are pair plasmas and there are observa-
tional indications that the composition nearer the AGN is dominated by
pairs (Wardle, et al. 1998). On the other hand it is difficult to create the
jets entirely with pairs, because of extreme energy losses, and some models
posit that the baryon load is dominant (e.g., Mannheim 1993; Falcke &
Biermann 1995).

While they are very powerful and may be capable of accelerating UHE-
CRs, radio jets are quite rare phenomena, especially in the local universe.
There are a few candidates inside the 100 Mpc range associated with the
GZK feature. Centaurus A is only about 3 Mpc away, for example, and
M87 is about 18 Mpc from us. Both of these are relatively low luminosity
objects, but might be adequate. Really high luminosity radio jets, such as
Cygnus A ($\sim 200$ Mpc) or quasars are generally much farther away. The
difficulty with sparse sources, as mentioned in §2, is how to produce an
apparently isotropic distribution of detected events. Early indications for
concentrations along the supergalactic plane, which might signal associa-
tion with matter concentrations, seem to have gone away with increased
data (Takeda, et al. 1999; Uchihori, et al. 2000). If the GZK feature is
seen, this may open the door to a larger scale distribution. The alterna-
tive possibility may still be open that intergalactic magnetic fields able to
redistribute the events.

5.2. GAMMA RAY BURSTS

Gamma Ray Bursts are an attractive idea for UHECRs, because they in-
volve highly relativistic flows and possibly very strong shocks. The more
common long-duration burst events are thought now probably to repre-
sent the formation of a black hole during core collapse of a massive star
ACCELERATION OF UHECRS

(Mészáros 2002), with evidence pointing towards the generation of relativistic jets, perhaps with Lorentz factors of several hundred. These high Lorentz factors can result when very little mass is incorporated ($< 10^{-5} M_{\text{sun}}$), and when the total kinetic energy in the flows is comparable to that in a conventional supernova, ($\sim 10^{51} \text{erg}$). The ejected plasmas (“fireballs”) are thought to be pair dominated, because they are so incredibly hot, although there will be a minor baryonic “load.” In this scenario variations in the energy ejection are supposed to lead to internal, mildly relativistic, shocks within the fireball, and those are plausibly capable of accelerating protons to UHECR energies (e.g., Waxman 2000). We looked at the energy constraints on CRs being accelerated in §3 and found that energies near the UHECR range were barely possible, with the most severe constraint being set by synchrotron losses in the posited magnetic fields. As noted in §2 if they were distributed uniformly in the universe, GRBs might be common enough to account for the measured flux of events. However, as pointed out by Stecker (2000) GRBs seem to be relatively rare today; they seem instead to have been much more common in the past. In that case, their local space density is much to small, according to Stecker, to account for UHECRs.

5.3. PULSAR WINDS

I argued before that UHECRs are probably extragalactic, and this is true. On the other hand, if they could be made of iron rather than protons, then it is perhaps still plausible to discuss galactic scenarios. The obvious immediate worry is how to explain the absence of an associated asymmetry in the distribution of detected events.

While protons near 1 ZeV are not significantly deflected in their motions by the galactic magnetic field, iron nuclei could still be influenced. At energies below 100 EeV the distribution of events might appear roughly isotropic, even if the sources were near the plane. In that spirit I will briefly outline a recent model for UHECRs based on young pulsars.

The magnetic fields typically associated with pulsars ($B \gtrsim 10^{12} \text{G}$) in concert with their rapid rotation ($\omega \lesssim 10^{3} \text{s}^{-1}$) could in principle accelerate protons to energies higher than 10 EeV, and a factor $Z = 26$ higher for Fe, according to the Hillas constraint. Magnetar fields are stronger, although the rotation periods are slower, so the nominal limits are similar. So, in that regard neutron stars are natural candidates for UHECRs. Venkatesan, et al. (1997) pointed out, however, that pair production should limit the maximum energy in the magnetosphere to about 1000 TeV, and that synchrotron losses near the neutron star are also severe. Blasi, et al. (2000) have argued, on the other hand, that relativistic pulsar winds outside the light circle are capable of accelerating Fe to energies above 100 EeV. This
is allowed by the Hillas constraint and also consistent with the energy per particle that would be carried away in a wind extracting the rotational energy of a rapidly rotating pulsar. Being outside the intense photon field and magnetic fields of the near magnetosphere, these particles should be able to escape. These authors also estimate that the combined efficiency of acceleration and galactic trapping need be only $\sim 10^{-5}$ to account for the observed flux, given a birthrate of pulsars in the galaxy around $10^{-2}$yr$^{-1}$.

Especially at the highest energies it is still difficult to produce a truly isotropic distribution of detected events in such a scenario. So, once again, the improved statistic expected in the near future from Auger should determine if any such class of models is really viable.

6. Conclusions

I have tried to outline in a simple fashion some of the underlying issues associated with the physics of accelerating charged particles to energies close to 1 ZeV. The evidence is strong that nature somehow produces such particles, either by acceleration or by decay from something even more energetic. They are rare, but their existence demonstrates something new. Both possible scenarios stretch our current understandings. Here I have taken the more conventional road supposing that physical environments exist where particles can reach such energies from below. There are few accelerators we have been able to imagine that apply sufficient energy gains and contain particles well enough for them to reach these energies. Even if that seems possible in principle, the energies losses faced by such particles are huge, due especially to collisions with ambient photons or by synchrotron radiation.

The bottom line is that very few candidate environments exist. Mostly they are extragalactic, both because it is difficult to confine the UHECRs in the galaxy, and because there are few galactic candidate phenomena. New proposals appear routinely, and some other suggestions may turn out to be obvious. At the moment, however, it seems to me that the strongest candidate accelerators are the high powered jets associated with radio galaxies, although they are not sure bets. These jets start out as moderately relativistic, they can carry their energy to large distances and they have lifespans of $10^8$ years or so. There are several plausible dissipation mechanisms that have been identified in radio jets that may be capable of producing very high energy particles, including shock acceleration, shear acceleration and even magnetic reconnection. We know from observations of nonthermal X-rays in some such objects that electrons have been generated with energies in the TeV range. UHECRs are probably hadronic, not leptonic, and that may be an issue. There are arguments favoring hadronic jets in these ob-
jects, although the more common view is that they are leptonic.

No matter where this topic goes it is certain to be interesting in the coming few years. Experiments under construction, such as the Auger Observatory, and those under discussion for the future, such as OWL, will vastly increase detection capabilities. Then we should finally know with confidence what role is played by photopion losses in UHECR propagation and how the detected events are distributed on the sky. That information should tell us much more clearly how the sources must be distributed, what their energy spectra must be and perhaps, if we are lucky, something about the propagation of the CRs, and thus information about intergalactic magnetic fields. Along the way we should learn finally if UHECRs are probes of fundamental physics or astrophysics.

Acknowledgements

I am most grateful to Professor Norma Sánchez for her invitation to participate in the delightful and stimulating 9th Chalonge School on Astrophysical Physics. My work on the physics of particle acceleration is supported by the NSF through grant AST00-71167, by NASA through grant NAG5-10774 and by the University of Minnesota Supercomputing Institute.

References

Abu-Zayyad, T., et al. (2002), astro-ph/0208301.
Achterberg, A., Gallant, Y. A., Kirk, J. & Guthmann, A. W. (2001), 328, 393.
Axford, W. I., Leer, E. & Skadron, G. (1977), Proc. 15th ICRC (Plovdiv), 11, 132.
Bell, A. R. (1978), Mon. Not. R. Astron. S., 182, 147.
Berezinsky, V., Gazizov, A. Z. & Grigorieva, S. I. (2001), astro-ph/0107306.
Biermann, P. L., Ahn, E.-J., Kronberg, P., Medina-Tanco, G. & Stanev, T. (2001), “Physics and Astrophysics of Ultra-High-Energy Cosmic Rays,” ed: M. Lemoine & G. Sigl, Lect. Notes Phys., 576, 181.
Bird, D. J., et al. (1994), Astrophys. J., 441, 144.
Blandford, R. D. & Ostriker, J. P. (1978), Astrophys. J., 221, L29.
Blasi, P, Epstein, R. I., & Olinto, A. (2000), Astrophys. J., 533, L123.
Erlykin, A. D., Lagutin, A. A. & Woffendale, A. W. (2002), astro-ph/0209506.
Falcke, H., & Biermann, P. L. (1995), Astron. Astrophys., 298, 395.
Fermi, E. (1949), Phys. Rev., 75, 1169.
Greisen, K. (1966), Phys. Rev. Lett., 16, 748.
Hayashida, N., et al. (1999), Astropart. Phys., 10, 303.
Hillas, A. M. (1984), Ann. Rev. Astron. Astrophys., 22, 425.
Jokipii, J. R., Kota, J. & Morfill, G. (1989), Astrophys. J., 345, L67.
Jones, T. W. (2001), in “7th Taipei Astrophysics Workshop on Cosmic Rays in the Universe,” ed: C.-M. Ko, ASP Conf. Ser., 241, 239.
Kang, H., Rachen, J. P., & Biermann, P. L. (1997), Mon. Not. R. Astron. Soc., 286, 257.
Kang, H., Ryu, D. & Jones (1996), Astrophys. J., 456, 422.
Kang, H., Jones, T. W. & Gieseler, U. D. J. (2002), Astrophys. J. (in press).
Kronberg, P. P. (2001), in “High Energy Gamma-Ray Astronomy,” ed: F. Aharonian & H. Völk, AIP, p 451.
Krymskii, G. F. (1977), Dokl. Akad. Nauk. SSSR, 234, 1306.
Lagage, P. O. & Cesarsky, C. J. (1983), Astron. Astrophys., 125, 249.
Lemoine, M., Sigl, G., Olinto, A. V., & Schramm, D. N. (1997), Astrophys. J., 486, 115.
Lesch, H. & Birk, G. T. (1998), Astrophys. J., 499, 167.
Malkov, M., Diamond, P. H. & Völk, H. J. (2000), Astrophys. J., 533, 171.
Malkov, M. & Drury, L. O'C. (2001), Rep. Prog. Phys., 64, 429.
Mannheim, K. (1993), Astron. Astrophys., 269, 67.
Medina Tanco, G. A., de Gouveia Dal Pino, E. M., & Horvath, J. E. (1997), Astropart. Phys., 6, 337.
Mészáros, P. (2002), Ann. Rev. Astron. Astrophys. 40, 137.
Mészáros, P. & Rees, M. J. (1994), Mon. Not. R. Astron. Soc., 269, L41.
Norman, C. A., Melrose, D. B. & Achterberg, A. (1995), Astrophys. J., 454, 60.
Olinto, A. V. (2000), Phys. Rept., 333-334, 329.
Ostrowski, M. (2000), Mon. Not. R. Astron. Soc., 312, 579.
Quest, K. (1988), in "Proc. 6th Intl. Solar Wind Conf., ed: T. O'Neil, & D. Brook (NCAR/TN-360), p 503.
Rachen, J. P. & Biermann, P. L. (1993), Astron. Astrophys., 272, 161.
Sigl, G. (2002), astro-ph/0210049.
Schopper, R., Birk, G. T. & Lesch, H. (2001), Proc. of ICRC 2001 (Hamburg), 2031.
Stecker, F. W. (2000), Astropart. Phys., 14, 207.
Stecker, F. W. (2002) (This volume) astro-ph/0208507.
Takeda, et al. (1999), Astrophys. J., 522, 225.
Takeda, M., et al. (2002), astro-ph/0209422.
Tregillis, I. L., Jones, T. W. & Ryu, D. (2001), Astrophys. J., 557, 475.
Uchihori, et al. (2000), Astropart. Phys., 13, 151.
Venkatesan, A., Miller, M. C. & Olinto, A. V. (1997), Astrophys. J., 484, 323.
Wardle, J. F. C., Homan, D. C., Ojha, R. & Roberts, D. H. (1998), Nature, 395, 457.
Watson, A. A. (2000), Phys. Rep., 333-334, 309.
Waxman, E. (2000), Astrophys. J. (Supplements), 127, 519.
Zatsepin, G. T. & Kuz'min, V. A. (1966), Zh. Esks. Teor. Fiz., Pis'ma Red., 4, 144.