ULTRA–HIGH-ENERGY COSMIC RAYS FROM HYPOTHETICAL QUARK NOVAE

R. OUYED
Department of Physics and Astronomy, University of Calgary, 2500 University Drive NW, Calgary, AB T2N 1N4, Canada; ouyed@phas.ucalgary.ca

P. KERÄNEN
Nordic Institute for Theoretical Physics, Blegdamsvej 17, DK-2100 Copenhagen, Denmark; and Centre for Underground Physics in Pyhäșalmi, P.O. Box 22, FIN-86801 Pyhäșalmi, Finland

AND

J. MAALAMPI
Department of Physics, P.O. Box 35, FIN-40014, University of Jyväskylä; and Helsinki Institute of Physics, P.O. Box 64, FIN-00014, University of Helsinki, Finland

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ABSTRACT

We explore the acceleration of ions in the quark nova (QN) scenario, in which a neutron star experiences an explosive phase transition into a quark star (born in the propeller regime). In this picture, two cosmic-ray components are isolated: one related to the randomized pulsar wind and the other to the propelled wind, both boosted by the ultrarelativistic QN shock. The latter component acquires energies $10^{15}$ eV $< E < 10^{18}$ eV, while the former, boosted pulsar wind, achieves ultrahigh energies $E > 10^{18.6}$ eV. The composition is dominated by ions present in the pulsar wind in the energy range above $10^{18.6}$ eV, while at energies below $10^{18}$ eV the propelled ejecta, consisting of the fallback neutron star crust material from the explosion, is the dominant one. Added to these two components, the propeller injects relativistic particles with Lorentz factors $\Gamma_{\text{prop}} \sim 1$–1000, later to be accelerated by galactic supernova shocks. The QN model appears to be able to account for the extragalactic cosmic rays above the ankle shown in the figure in this paper and to contribute a few percent of the galactic cosmic rays below the ankle. We predict a few hundred ultra–high-energy cosmic-ray events above $10^{19}$ eV for the Pierre Auger detector per distant QN, while some thousands are predicted for the proposed Extreme Universe Space Observatory (EUSO) and Orbiting Wide-angle Light collectors (OWL) detectors.

Subject headings: acceleration of particles — cosmic rays — elementary particles

1. INTRODUCTION

Large efforts have been devoted to explore the origin of cosmic rays. Most puzzling are the observed ultra–high-energy cosmic-ray events (UHECRs) above the Greisen-Zatsepin-Kuzmin (GZK) cutoff (Greisen 1966; Zatsepin & Kuzmin 1966): protons lose energy drastically due to pion photoproduction processes on the cosmic microwave background (CMB) at energies higher than $7 \times 10^{19}$ eV, and this limits the proton mean free path to some tens of megaparsecs (see Blasi [2003] and Bahcall & Waxman [2003] for more discussion on the GZK cutoff). Two distinctly different classes of models are commonly considered: the top-down and bottom-up scenarios. In the former, UHECRs are associated, e.g., with the decay of some supermassive particles, whereas in the latter, UHECRs are assumed to be accelerated by astrophysical objects (e.g., Sigl et al. 1994; Venkatesan et al. 1997; Biermann 1998; for a more recent review we refer the interested reader to Ostrowski 2002, Arons 2003, Sigl 2003, and Olinto 2004). A growing number of bottom-up explanations have been proposed, including active galactic nuclei, gamma-ray bursts, and neutron stars. However, at present there seems to be no clear association of UHECR events with any of these objects, although by tuning the model parameters one enables the highest energy events to be accounted for.

Identifying sources of the observed UHECR events remains uncertain and debated, leaving room for speculation (Nagano & Watson 2000; Torres & Anchordoqui 2004; Staney 2004). As such, we wish to consider in this paper the possibility that hypothetical quark novae (QNe; Ouyed et al. 2002) contribute to the cosmic-ray flux, especially above $10^{18}$ eV. In the QN explosion the core of a neutron star (NS) shrinks into the equilibrated quark object/star (QS). The overlaying crust material free-falls following the core contraction, releasing enough energy to form an ultrarelativistic ejecta. The ultrarelativistic shock interaction with its surroundings environment (namely, the randomized relativistic wind of the progenitor) leads to the first cosmic-ray component in our model. The compact remnant acting as a magnetohydrodynamic propeller provides the second component, injecting ions to be boosted by the ultrarelativistic QN shock. Most of the propelled wind of relativistic particles is injected into interstellar space with low energies (from GeV to TeV), and these particles can be further accelerated by galactic supernova shocks. Thus, the propeller can contribute to the cosmic-ray flux also below the knee shown in Figure 1 as an injector. It appears that the QS model can largely contribute to the extragalactic cosmic rays above the ankle and contribute a few percent of the galactic cosmic rays around and below the knee.

The main assumption in this paper is the conversion of a highly magnetized, rapidly spinning NS to a QS. This phase conversion (or second explosion), which remains to be confirmed, leads to unique conditions in which two UHECR components are feasible. It is this unique aspect of the model that makes this investigation worth pursuing. We note, however, that the approach presented here (including the acceleration mechanisms as related to the dynamics of the explosion) borrows heavily from what has already been presented in the literature for models of UHECRs.
in the context of isolated pulsars and binary coalescence (e.g., Venkatesan et al. 1997, Gallant & Achterberg 1999, Blasi et al. 2000, and Arons 2003 to cite only a few).

This paper is presented as follows. We start with a brief review of the concept of a QN and its features (§ 2). In § 3 we describe conditions under which the compact remnant is born in a propeller regime. In § 4 we explain how cosmic rays are produced and isolate the two components (with energies $10^{15} \text{ eV} < E < 10^{18} \text{ eV}$ and $E > 10^{18.6} \text{ eV}$). Here we describe the QN compact remnant as an injector of relativistic particles to the galaxy. In § 5 we predict UHECR events in future detectors and discuss multiple events and the isotropy of sources in our model. A discussion and a conclusion follow in § 6.

2. QUARK NOVA

It has been suggested that the core of a NS may deconfine to a composition of up ($u$) and down ($d$) quarks during or shortly after some supernova explosions during which the central density of the proto-NS is high enough to induce phase conversion (see, e.g., Dai & Peng 1995; Xu et al. 2001). It has also been speculated that when the density in the core increases further, a phase with strange quarks ($s$) becomes energetically favored over the pure ($u,d$) phase, and soon the entire star is contaminated and converted into this ($u,d,s$) phase. This is one of the scenarios introduced to convert an entire NS to a QS (e.g., Cheng & Dai 1996; Bombaci & Datta 2000). In the QN picture the ($u,d,s$) core is assumed to shrink to a corresponding stable QS before the contamination is spread over the entire star. By physically separating from the overlying material (hadronic envelope) the core drives the collapse (free fall) of the left-out matter, leading to both gravitational energy and phase-transition energy release as high as $E_{\text{QN}} \approx 10^{53} \text{ ergs}$ (Ouyed et al. 2002; Keränen et al. 2005).

2.1. Quark Nova Ejecta and Compact Remnant

The QN ejecta consist mainly of heavy nuclei (the NS crust), protons, electrons, and neutrons. The total amount of the ejecta has been estimated to be of the order of $(0.001\sim0.01) M_\odot$ with corresponding Lorentz factor $\Gamma_{\text{QN}} = \epsilon E_{\text{QN}}/m_{\text{QNe}}c^2 \approx 10\sim100$, where $\epsilon$ is the percentage of the QN energy transferred to the QN shock front, which we assume to be $10\%$.

The radius of the newly formed QS is given by $R \approx R_{\text{NS}}(\rho_{\text{NS}}/\rho)^{1/3}$, where the NS core density, $\rho_{\text{NS}}$, and the QS density, $\rho$, scale as $\rho_{\text{NS}}/\rho \approx 0.1\sim0.2$ (Ouyed et al. 2002; Keränen et al. 2005); hereafter, parameters with no subscripts refer to the QS. The QS spins up during the phase transition because of contraction, with a rotational period $P = P_{\text{NS}}(R_{\text{NS}}/R)^3 P_{\text{NS}}(\rho_{\text{NS}}/\rho)^{2.3} \approx P_{\text{NS}}/4$. In the process, the magnetic field is amplified, $B = B_{\text{NS}}(R_{\text{NS}}/R)^2 = B_{\text{NS}}(\rho/\rho_{\text{NS}})^{2/3} \approx 4B_{\text{NS}}$.

2.2. Fallback Material

Fallback of material onto a newly formed QS may take place (similar to what has been suggested in the supernova case; Chevalier 1989). In the early stages the accretion rate is given by

$$\dot{m} \approx 10^{28} \text{ g s}^{-1} \left(\frac{\rho_{\text{fb}}}{10^9 \text{ g cm}^{-3}}\right) \left(\frac{R}{10 \text{ km}}\right)^{3/2} \left(\frac{M}{1.5 M_\odot}\right)^{1/2},$$

(1)

where $\rho_{\text{fb}}$ is the average density of fallback matter ($10^6 \text{ g cm}^{-3}$, representing the crust material and the matter below the neutron drip line densities). The hyper-Eddington accretion rate given above is understood by noting that (1) the fallback matter in the initial phase of the explosion is not accreted onto the surface (but is propelled away before radiating) and (2) the crust would not form in the early stages leaving the QS bare (not subject to the Eddington limit, since the bulk of the star is bound via strong interaction rather than gravity; see, e.g., Alcock et al. [1986], Zhang et al. [2000], and Xu [2003] for a recent discussion on the matter).

3. PROPELLER REGIME

For the remainder of this paper, we consider a QS (NS) mass of $M = 1.5 M_\odot$ ($M_{\text{NS}} \approx 1.5 M_\odot$), a radius $R = 10 \text{ km}$ ($R_{\text{NS}} \approx 12.5 \text{ km}$), and a surface magnetic field of $4 \times 10^{14} \text{ G}$ ($B_{\text{NS}} \approx 10^{14} \text{ G}$). Since the fastest pulsar has a period of 1.56 ms (Backer et al. 1982), we use 2 ms as a representative value for the period $P$ of the newborn QS ($P_{\text{NS}} \approx 8 \text{ ms}$). Clearly these are unique constraints (millisecond, highly magnetized NS) favoring a scenario in which the QS forms immediately following a supernova explosion or in which the NS has been spun up by accretion from a companion.

The newly born QS is defined by its three critical radii: the Keplerian “corotation radius”

$$R_{\text{cor}} = 27 \text{ km} \left(\frac{M}{1.5 M_\odot}\right)^{1/3} \left(\frac{P}{2 \text{ ms}}\right)^{2/3},$$

(2)

the magnetospheric radius at which the ram pressure of the infalling matter balances the magnetic pressure (see, e.g., Frank et al. 1992)

$$R_{\text{mag}} = \left(\frac{B^2 R^6}{2 \mu} \sqrt{2GM}\right)^{2/7} = 55 \text{ km} \left(\frac{B}{4 \times 10^{14} \text{ G}}\right)^{4/7} \times \left(\frac{10^{28} \text{ g s}^{-1}}{m}\right)^{2/7} \left(\frac{R}{10 \text{ km}}\right)^{12/7} \left(\frac{1.5 M_\odot}{M}\right)^{1/7},$$

(3)
and the light cylinder radius

\[ R_{lc} = \frac{c}{\Omega} = 96 \text{ km} \left( \frac{P}{2 \text{ ms}} \right). \]  

(4)

Given our fiducial values, the QS is born in the propeller regime (Shvartsman 1970; Illarionov & Sunyaev 1975), i.e., \( R_{\text{crit}} < R_{\text{mag}} < R_{lc} \), in which the infalling material may be accelerated in a wind that carries away angular momentum from the magnetosphere and hence from the QS itself (\S 4.2). Note that the propeller regime is only possible if the QS is born with a magnetic field exceeding the critical value \( B_{c} \geq 10^{14} \, \text{G} \), which corresponds to a parent NS initial magnetic field \( B_{\text{NS,c}} = B_{c}/4 \geq 2.5 \times 10^{13} \, \text{G} \). Here “\( \text{c} \)” stands for “critical,” while the factor 4 (derived from our fiducial values) is the result of the amplification of the magnetic field following the QN explosion. These critical values explain our assumption of a highly magnetized, rapidly spinning NS.

### 3.1. Propeller Lifetime

The star loses rotational energy at a rate defined by the gravitational radiation losses \( \dot{E}_{\text{grav}} \), electromagnetic radiation losses \( \dot{E}_{\text{em}} \), and the propeller’s torque \( \dot{E}_{\text{prop}} \). That is,

\[ \frac{d\Omega}{dt} = -\dot{E}_{\text{grav}} - \dot{E}_{\text{em}} - \dot{E}_{\text{prop}} / I \Omega \]  

(5)

with (Shapiro & Teukolsky 1983)

\[ -\dot{E}_{\text{grav}} = \frac{9}{5} \frac{G I^{2} \lambda^{2} \Omega^{6}}{c^{5}} = 6.8 \times 10^{47} \, \text{ergs/s} \]

\[ \times \left( \frac{2 \text{ ms}}{P} \right)^{6} \left( \frac{I}{10^{45} \, \text{g cm}^{2}} \right)^{2} \left( \frac{\epsilon}{0.01} \right)^{2}, \]

(6)

where \( I \) is the moment of inertia and \( \epsilon \) the equatorial eccentricity, and with (Manchester & Taylor 1977)

\[ -\dot{E}_{\text{em}} = \frac{4B^{2}R^{6} \Omega^{4}}{9c^{3}} \times 2.6 \times 10^{46} \, \text{ergs/s} \]

\[ \times \left( \frac{2 \text{ ms}}{P} \right)^{4} \left( \frac{B}{4 \times 10^{14} \, \text{G}} \right)^{2} \left( \frac{R}{10 \, \text{km}} \right)^{6}. \]

(7)

The spin-down due to the propeller is (Menou et al. 1999)

\[ -\dot{E}_{\text{prop}} = 2m \dot{c} e^{2} = 1.8 \times 10^{49} \, \text{ergs/s} \]

\[ \left( \frac{\dot{m}}{10^{28} \, \text{g s}^{-1}} \right). \]

(8)

The expression above assumes that the material flung away by the propeller effect has been accelerated to an angular speed corresponding to that of the star. We note that gravitational losses would be important in the spin-down for very short periods \( P < 2 \, \text{ms} \) (not our case), and propeller losses dominate over the electromagnetic dipole radiation losses given our accretion rates.

The propeller regime lifetime can be obtained using equations (5) and (8). We find, assuming a constant accretion rate,

\[ t_{\text{prop}} \simeq 10^{3} \left( \frac{I}{10^{45} \, \text{g cm}^{2}} \right) \left( \frac{10^{28} \, \text{g s}^{-1}}{\dot{m}} \right) \left( \frac{1}{P_{i}} - \frac{1}{P_{f}} \right). \]

(9)

where \( P_{i} \) and \( P_{f} \), initial and final periods, are in milliseconds. For a constant accretion rate, and given the total amount of the QN ejecta, the lifetime of the propeller phase would not exceed 100 s, during which time the QS would have spun down by no more than 30%, thus remaining within a period of milliseconds. Following this phase, the spin-down is governed by the magnetic dipole radiation losses of the QS.

### 4. PRODUCTION OF THE TWO COSMIC-RAY COMPONENTS

#### 4.1. Acceleration of the Pulsar Wind

We assume that the pulsar wind is composed of electron-positron pairs and ions, as would be expected if pulsars operate as open-circuited systems with ions carrying the return current (Hoshino et al. 1992). Some fraction of ions can be heavy like iron, originating from the surface of the pulsar. We also assume that the wind has been subject to a termination shock and has acquired a randomized direction of particle momenta (Gallant & Achterberg 1999). The number density of electron-positron pairs is \( n_{\text{e}+\text{p}} = B_{\text{NS}}(r)/\Omega_{\text{NS}}/(4\pi ec) \), where \( n_{\text{e}+\text{p}} \) is the Goldreich-Julian density (Goldreich & Julian 1969) and \( e \) is the electron charge.

We adopt \( \gamma_{w} \sim 10^{6} \) as the relativistic factor of most of the wind particles (see, e.g., Gallant & Arons 1994) and the ratio of ions to electron-positron pairs to be \( \alpha = n_{i}/n_{\text{e}+\text{p}} \sim 10^{-3} \); ions carry most of the energy of the wind. For future purposes, we write the ion density as

\[ n_{i}(r) = \alpha n_{\text{e}+\text{p}} \left( \frac{R_{\text{NS}}}{r} \right)^{\beta}, \]

(10)

where the index \( \beta \) describes the radial dependency.

The QN ejecta will interact with the parent NS wind. The randomized wind particles will be boosted with one shock crossing by a factor of \( 2f_{\text{QN}}^{2} \sim 2 \times 10^{46} \) before they leave the region.\(^1\) The maximum number of particles that can be accelerated is limited by the QN fireball energy and can be written as (here the energy of the pairs will be negligible, since energetically they are at much lower gamma factors due to synchrotron losses)

\[ N_{w} \sim \frac{\epsilon E_{\text{QN}}}{2 \Gamma_{\text{QN}}^{2} \gamma_{w}^{2} Z m p c^{2}} \sim \frac{3\epsilon}{Z} \times 10^{45}, \]

(11)

where \( m_{p} \) is the proton mass, \( Z \) can vary from 1 to 26 (iron), and \( \epsilon \) is the fraction of QN energy transferred via the QN shock to the pulsar wind particles. If all ions are protons and approximately accelerated to the energies of the ankle (around \( 10^{18.5} \, \text{eV} \)), then up to \( 10^{45} \) particles would gain that energy. Heavier ions will correspondingly reach higher energies with the same total relativistic factor \( 2\Gamma_{\text{QN}}^{2} \gamma_{w} \), e.g., iron will easily gain energies over \( 10^{20} \, \text{eV} \). Therefore, in this model at higher energies there should be a natural increase in heavier ions. Since the above follows closely calculations already presented in Gallant & Achterberg (1999), the corresponding spectrum also in our case can be shown to be \( dN/dE \sim E^{-2} \) or even flatter given its dependency on \( \Gamma_{\text{QN}} \). We would like to emphasize that further studies of our model are needed in order to make better predictions of the spectrum.

Such an injection spectrum might agree with AGASA measurements, but let us recall that in the high-energy region we discuss here, currently there seems to be a disagreement between the AGASA ground array (Takeda et al. 1999) and the HiRes fluorescence detector (Abbasi et al. 2005), which seems consistent with...
with the GZK cutoff. Clearly there is a need for much larger experiments such as the Pierre Auger,\(^2\) Extreme Universe Space Observatory (EUSO), and Orbiting Wide-angle Light collectors (OWL) that can increase the number of detected events by 1 or 2 orders of magnitude before the injection spectrum is known conclusively. As we have said, for now, our model lacks the details to predict the exact spectrum.

The number of particles per unit volume from QNe in all the galaxies is \(J_{\text{QN}} = (c/4\pi)\frac{n_* n_{\text{loss}}(E_{\text{qs}})}{E_{\text{qs}}}\) where \(n_{\text{loss}}(E_{\text{qs}})\) is the residence time of a particle at the ankle initially injected at higher energy (Arons 2003). The QN rate per galaxy is given by \(v_{\text{QN}}\), while \(n_*\) is the galaxy density. With \(v_{\text{QN}} = 10^{-6}\) yr\(^{-1}\) and \(n_* = 0.02\) Mpc\(^{-3}\), this implies

\[
J_{\text{QN}}(E > 10^{18.8}\text{ eV}) \sim \frac{c}{Z} (5.31 \times 10^{-18}) \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1},
\]

(12)
as compared to the observed value \(J_{\text{obs}}(E > 10^{18.8}\text{ eV}) \sim 3 \times 10^{-10}\) cm\(^{-2}\) s\(^{-1}\) sr\(^{-1}\) (one event per square kilometer per year; Lawrence et al. 1991; Takeda et al. 1998), assuming \(Z = 1\). For particles at energies higher than \(10^{18.8}\) eV we used an average \(v_{\text{loss}} \sim 1\) Gyr (Biermann & Strittmatter 1987; Berezinsky & Grigor’eva 1988; Bertone et al. 2002; see also Fig. 5 in Sigl 2004b).

The QN flux calculated above is short of the observed one unless the QN is very efficient in transferring energy via the QN shock into relativistic particles. However, the presence of a large-scale magnetic field complicates the interpretation of UHECR data and could lead to an overestimate of the observed flux. More specifically, the energy spectrum of particles emitted from a (nearby) source depends very strongly on the relative position of the observed source and the observer to the direction of the large-scale field. In some cases the flux can be enhanced by a factor of \(\sim 100\) for energies below \(10^{19}\) eV (see Stanev et al. 2003 and references therein). The observed flux and spectrum might not reflect the ones at the source, making uncertain a direct comparison with the numbers derived above.

### 4.1.1. Acceleration Timescale

The ejecta remain in expansion phase until the shock starts to decelerate. The transition between these two phases occurs at a radius \(R_{\text{Qn,d}}\) at which the energy in the swept-up material becomes of the order of the energy released in the QN, i.e., when a large fraction of the fireball energy has been used to reaccelerate the wind particles (Blandford & McKee 1976; Gallant & Achterberg 1999). Using equation (10),

\[
R_{\text{Qn,d}} \simeq \left[ \frac{E_{\text{Qn}}(3 - \beta)}{4\pi c \rho_{\text{IGR}} R_{\text{Qn}}^2 \tau_{\text{loss}} Z n_p c^2} \right]^{1/(3 - \beta)},
\]

(13)
which gives \(R_{\text{Qn,d}} \sim 5.5 \times 10^{13}/\sqrt{Z}\) cm for \(\beta = 1\). If the ion density decreases as \(\sim 1/r^2\) (\(\beta = 2\)), the fireball will be in the expansion phase much longer and the distance should be defined by using interstellar matter density together with the wind (we estimate \(R_{\text{Qn,d}} \geq 10^{15}\) cm for \(\beta = 2\)). These numbers will define the timescale of the highest energy cosmic-ray component acceleration. The case with \(\beta = 1\) would more closely follow the magnetic field value and therefore the charge density of the pairs. We use it as an example. The corresponding time is \(t_{\text{Qn,d}} = R_{\text{Qn,d}}/c \sim 1700/\sqrt{Z}\) s. If all wind particles are iron, then the timescale would be shorter, \(330\) s. One should notice, however, that the ion density profile in the pulsar wind bubble is not known, and therefore the timescale of the expansion phase could be much larger (with \(\beta \sim 2\) up to weeks).

### 4.2. Acceleration of the Propelled Wind

The second cosmic-ray component occurs when rapidly rotating millisecond pulsars with large magnetic fields, \(B > 10^{14}\) G, undergo a QN explosion. In this case the born QS is in the propeller regime (see § 3). The material flung away by the propeller expands as a magnetohydrodynamic wind to reach the speed of light (with Lorentz factors up to 1000 at the light cylinder for \(B > 10^{14}\) G; see Fendt & Ouyed [2004] for detailed calculations).\(^3\) The magnetohydrodynamic nature of the propeller ensures that acceleration occurs at large distances from the star where synchrotron losses are likely to be minimal. As demonstrated in Fendt & Ouyed (2004), the magnetic field above the Alfvén surface is predominantly toroidal. Such a geometry will allow the particles to escape freely along the poloidal direction in the acceleration zone without being deflected by the magnetic field lines. Particles with \(\mathbf{v}_\text{prop} > \mathbf{v}_\text{QN}\) can reach eventually the QN shock in order to reach Lorentz factors as high as \(\mathbf{v}_\text{prop}/\mathbf{v}_\text{QN} \approx 2 \times 10^7\). It is roughly \(10^{16}\) eV for protons and \(10^{18}\) eV for iron in one shock crossing. Since an efficient propeller needs strong magnetic field and a short period, this component is created in young pulsars soon after the supernova explosion or in binaries in which the NS spins up (which also provides extra mass to trigger the QN).

If the propeller works with the extreme accretion rate of \(m = 10^{28}\) g s\(^{-1}\), then the rate of particle number is \(N_{\text{prop}} = m/\rho_{\text{IGR}} \sim 10^{52}/Z\text{ s}^{-1}\). That is, after one crossing, during which a particle is boosted by a factor \(2 \Gamma_{\text{QN}}\), the energy per unit time used to accelerate the propeller wind is \(N_{\text{prop}} m_{\text{p}} c^2/2\); only the particles with \(\mathbf{v}_\text{prop} \geq \mathbf{v}_\text{QN}\) will be accelerated by the QN shock. This indicates that particle acceleration consumes the energy of the shock wave in \(\sim 1\) ms, and the shock dies out.\(^4\) We expect the propeller to function efficiently in its early stages with most particles acquiring \(\mathbf{v}_\text{prop} \geq \mathbf{v}_\text{QN}\). In reality this timescale is longer because the propeller wind particles cannot propagate straight forwardly to the shock (due to magnetic field and turbulence); their energies must have spread out and the most energetic particles reach the shock first. Nevertheless, it is short compared to the timescale of the expansion phase. It also means that a tiny amount of the bulk of the propeller wind will be accelerated to energies of the order of \(10^{15}\) eV and at most up to \(10^{18}\) eV for those few particles that managed to get one more kick in the process. Here again, given the many similarities to the calculations and approach presented in Achterberg et al. (2001), we expect the boosted propeller wind to acquire a power-law spectrum, \(dN/dE \sim E^{-\gamma}\) with \(\gamma \approx 3.2\)–3.3.

Using the Milky Way dimensions with a radius of 15 kpc and a scale height of \(\sim 1\) kpc, we obtain the number density of galactic cosmic rays for one QN to be of the order of \(n_{\text{MW}} \sim 4 \times 10^{-18}\) cm\(^{-3}\) \(Z(10^{15}\text{ eV}/E) E_{\text{Qn}}/(10^{60}\text{ ergs})\). We assume a leaky box model (Garcia-Munoz et al. 1987; Simpson & Garcia-Munoz 1987).

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\(^2\) For more information about the Pierre Auger Observatory, see http://www.auger.org.

\(^3\) Shvartsman (1970) first suggested that relativistic particles can be formed at the propeller stage by a rapidly rotating magnetic field (see also Kundt 1990).

\(^4\) Some of the ejecta may fall back, allowing for disk and later planet formation around the newly born QS (Keränen & Ouyed 2003).
1988) for the galactic cosmic rays diffusing out of the galaxy in which we adopt the leakage timescale expressed as \( T_d = 3 \times 10^7 \text{ yr} / (E_{\text{Gev}} / Z)^{-1/3} \), where \( T_d \) is diffusion time (see, e.g., Webber et al. 1998; Biermann et al. 2001; Biermann & Sigl 2001). The estimated time-averaged flux in our model is

\[
J^{\text{QN}}(E > 10^{15} \text{ eV}) \sim \frac{c}{4\pi} n_{\text{MW}} T_d \nu_{\text{QN}} \\
\approx \frac{3 \times 10^{-10}}{Z^{2/3}} \left( \frac{10^{15} \text{ eV}}{E} \right)^{4/3} \left( \frac{\nu_{\text{QN}}}{10^{-6} \text{ yr}^{-1}} \right) \\
\times \left( \frac{\epsilon E_{\text{QN}}}{10^{52} \text{ ergs}} \right) \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1},
\]

which is similar to the observed value of \( J^{\text{obs}}(E > 10^{15} \text{ eV}) \sim 3 \times 10^{-10} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \) (Bird et al. 1995; Candia et al. 2003; Haungs et al. 2004) as long as the hyper-Eddington accretion rates can be accepted and if the accelerated particles are mainly protons. Note that below the knee the \( J^{\text{QN}} \) is too low to account for the observed cosmic-ray flux; the shock is not energetic enough to accelerate more particles as to extend the spectrum to lower energies. Therefore, the steeper spectrum above the knee provided by the QN shock cannot dominate at energies below the knee. More important, the flux in equation (14) is averaged over time, \( t \gg 1/\nu_{\text{QN}} \). Above the knee the cosmic-ray diffusion time is shorter than the QN occurrence, \( T_d < \tau_{\text{QN}} = 1/\nu_{\text{QN}} \). For example, at the knee \( T_d \approx 3 \times 10^8 \text{ yr} \) for protons, and it is \( T_d \approx 3 \times 10^4 \text{ yr} \) at the ankle, while the QN occurrence time is \( \tau_{\text{QN}} \approx 10^6 \text{ yr} \). Therefore, the propelled wind of the QN model can account for the galactic cosmic-ray flux locally, i.e., in the vicinity and shortly after the QN event. Over the whole galaxy and in timescales longer than millions of years, QNe can contribute on average a few percent of the galactic cosmic rays around the knee. Closer to the ankle the contribution is negligible because of the very short diffusion time \( T_d \). More accurate estimates of the leakage times at high energies and QNe occurrence would be needed to make the above conclusions firm. In particular, the QN rate carries substantial uncertainties, among them the difficulty of determining the critical density for the transition to quark matter, the burden of not knowing which equation of state better describes the parent NS, and the lack of a complete theory that can describe the QCD phase diagram that would describe the path followed by the NS during its transition to a QS. We note, however, that recent studies show that QN rates can be as high as \( 10^{-4} \text{ yr}^{-1} \) \( \text{galaxy}^{-1} \) (Yasutake et al. 2004), which is above our fiducial value of \( 10^{-6} \text{ yr}^{-1} \) \( \text{galaxy}^{-1} \).

Finally, since the propelled wind consists of the crust material of the parent NS, we expect this cosmic-ray component to be rich in heavier ions, up to iron. Their rigidity \( (E/Z) \) means that these heavy nuclei would leak out of the galaxy at a lower rate than protons. It implies an increase of the heavy nuclei to hydrogen ratio in the chemical composition over time for those cosmic rays originating in QNe.

4.3. Quark Nova as a Relativistic Particle Injector

In our model the propeller injects relativistic particles. As discussed above, some of this propelled material can be reaccelerated by the QN shock. The remaining bulk will populate interstellar space. These particles with Lorentz factors \( \Gamma \sim 1-1000 \) (or energies \( 1-1000 \text{ GeV} \) for protons) could later be reaccelerated by galactic supernova shocks. For the number of injected particles a simple estimate gives \( N_{\text{inj}} \sim E_{\text{rot}} / (10 \text{ GeV}) \sim 10^{52} / 10^{-2} \text{ ergs} \sim 10^{54} \) that are accelerated to GeV energies by the propeller. Here \( E_{\text{rot}} \) is the maximum rotational energy of the QS. Using equation (14) this implies \( J \sim 10^{-3} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \), which is a few percent of the observed value (see, e.g., Abe et al. 2003). In this case the particle density in the galaxy can reach \( n_{\text{MW}} = 4 \times 10^{-13} \text{ cm}^{-3} \) from a single QN using up the QS rotational energy. The spectrum of the population of these injected particles will be shaped by later encounters and acceleration by galactic supernova shocks. We note here as well that at most a few percent of the cosmic rays below the knee can be of the QN origin.

The QN model predicts three cosmic-ray components, two galactic ones, and an extragalactic one. This should, in principle, lead into transitions or steps in the cosmic-ray flux. The galactic components will be subject to many supernova shocks that will smoothen the discontinuity around the knee. Note however that these signatures will be dwarfed by the much higher contribution from other galactic cosmic-ray sources. Having shown that QNe are candidate sources of UHECRs above the ankle, we next discuss how observational features can be explained within the QN model.

5. UHECRS FROM QUARK NOVAE IN FUTURE COSMIC-RAY DETECTORS

With a rate of \( 10^{-6} \text{ galaxy}^{-1} \text{ yr}^{-1} \) and a galaxy density \( n_g \approx 0.02 \text{ Mpc}^{-3} \) one should expect about 0.08 QNe each year within a 100 Mpc sphere. This 100 Mpc region contains galactic and intergalactic magnetic field regions to which GZK-energy cosmic rays are subject. This may lead to an arrival-time delay and therefore a clustering of events around the source. The arrival timescales can spread from years to millions of years depending on the strength and the configuration of the magnetic field as well as the distance to the QN. Here we refer the interested reader to Sigl (2004a) for a recent discussion on the effect of the magnetic field on the UHECRs. Given these numbers, there may be a tiny energy window with the particles clustering within a few degrees toward the source (e.g., Kronberg 1994a, 1994b). Such events, detected by future detectors, will signal the beginning of the era of observational particle astronomy. Naturally, such observations will be better guided if the QN can be detected by other means as well (e.g., X-rays, gamma rays).

The energy lost by the UHECRs as they propagate and interact with the cosmic microwave background is transformed by cascading into secondary GeV–TeV photons (Protheroe & Stanev 1996). This TeV–gamma-ray–UHECR trace could, in principle, be detected in the future (Catanese & Weekes 1999) and be used to test our model (see the discussion in Akerlof et al. 2003 and references therein). However, photons and UHECRs may have very different arrival times that are not easily quantified. This again calls for better understanding of magnetic field effects before such connections between the TeV photons and the related UHECRs can be made. In addition, since protons can be accelerated up to \( 10^{21} \text{ eV} \) in our model, significant neutrino fluxes (with energies above \( 10^{19} \text{ eV} \)) can be generated (e.g., Waxman & Bahcall 1999; Engel et al. 2001). This is below the currently advertised threshold of \( 5 \times 10^{19} \text{ eV} \) for EUSO and OWL, and most of the potential events will go undetected (Halzen & Hooper 2002). Nevertheless, future neutrino detectors should be able to signal any neutrino-UHECR trace with the arrival direction of ultra–high-energy neutrinos as a good indication of the QN location. The above-mentioned traces are the subject of another study of QNe as sources of UHECRs.

5.1. Multiple Events and Clustering of UHECRs

Assuming that the cosmic rays above the ankle are accelerated in QNe, we obtain roughly \( 10^{44} \) particles above
$10^{19}$ eV per QN. The integrated flux $F$ per unit area from one QN is

$$F \approx 10^{-10} \left( \frac{100 \text{ Mpc}}{D_{\text{QN}}} \right)^2 \text{cm}^{-2},$$

where $D_{\text{QN}}$ is the QN distance. This is the total flux integrated over time including the spread in arrival time induced by magnetic fields, which can vary from years to millions of years (Sigl 2004a). In a detector like Auger this implies (integrated over the year and aperture and assuming that the QN is within the sky coverage) a few hundred events per QN at a distance of 100 Mpc. In a detector like EUSO and OWL one can expect 1–2 orders of magnitude more events per QN. If the time delay and thus the spread in arrival time is, e.g., 1000 yr, one could expect doublets and triplets around a given QN within 10 yr in Auger and, equivalently, dozens of clustered events with EUSO and OWL. For completeness, if the average time delay of cosmic rays above $10^{19}$ eV is around 1000 yr, in our model it would mean that UHECRs from roughly 100 QNe can presently be detected within 100 Mpc.

5.2. Isotropy of UHECRs

Unlike the arrival directions of UHECRs, the galaxy distribution is not isotropic within 100 Mpc. This may mean that UHECR accelerators are not all within galaxies. It is natural to attribute UHECR sources to galaxies and the almost perfect isotropy, so far observed both below and above the ankle (e.g., Torres & Anchordoqui 2004), to magnetic fields. Despite the uncertainties on the magnetic field strength, it has indeed been shown that if our Local Supercluster contains a large-scale magnetic field it can provide sufficient bending to the cosmic-ray trajectories (Sigl et al. 1999; Farrar & Piran 2000). Heavy ions like iron, if present in the pulsar wind, will have unique consequences. Iron has a higher cutoff energy (Puget et al. 1976; Stecker & Salomon 1999) than lighter nuclei, thus allowing a longer path of propagation in a tiny energy window allowing for more distant sources to contribute to the flux. Furthermore, being bent in galactic and intergalactic fields, heavy nuclei would acquire a more isotropic distribution of arrival directions than protons. Finally, heavy nuclei will naturally acquire higher energies than protons and contribute to the extreme end of the UHECR flux in our model.

Of interest to the QN model, old population pulsars with large peculiar velocities could cover tens of megaparsecs from their origins in a Hubble time. Some of these runaway pulsars may undergo a QN explosion as a result of an increase of their core densities following spin-down or accretion from the surrounding space. This would lean toward a more isotropic distribution (the extent of isotropy in the runaway pulsar population remains to be determined) of the arrival directions of the UHECRs (the QN shock boosted pulsar wind). This favors weak magnetic field ($B < 10^{10}$ G) millisecond pulsars in our model, with the assumption that these old, isolated pulsars still sustain their wind bubbles. We note that these candidates are probably not born in the propeller regime following the QN explosion, given their weak magnetic field and long periods (see § 3).

5 It is a well-known fact that pulsars in our galaxy have velocities much in excess of those ordinary stars (Harrison et al. 1993). It is reported that their transverse speeds range from 0 to $\sim 1500$ km s$^{-1}$ and their mean three-dimensional speed is $450 \pm 90$ km s$^{-1}$ (Lyne & Lorimer 1994).

5.3. Injection Spectrum and Time Dependency

For a given distant extragalactic QN source the more energetic particles should arrive first, since the associated time delay induced by the intergalactic magnetic fields is short. Therefore, it would appear as if there is a monoenergetic flux of particles from the source at one given time. Later, the lower energy particles will arrive with a larger time delay (G. Sigl 2004, private communication).

In the case of galactic cosmic rays, the propelled wind consists of the crust material of the parent NS; the rigidity $(E/Z)$ implies that the heavy nuclei would leak out of the galaxy at a lower rate than that of protons. We thus expect an increase of the heavy nuclei to hydrogen ratio in the chemical composition over timescales of some ten thousand years for cosmic rays originating in QNe. Note, however, that this effect is local and could have importance only if we detect a nearby QS. If the star has undergone a QN phase in the past, the local cosmic-ray composition may reflect the time passed since the QN explosion.

6. SUMMARY AND CONCLUSION

The cosmic-ray acceleration in the QN model consists of three different components as illustrated in Figure 1. The first component is due to the acceleration of the high-energy particles in the pulsar wind bubble by the QN shock. We predict a few hundred UHECR events above $10^{19}$ eV for a distant QN with the Pierre Auger detector, while some thousands are predicted with the proposed EUSO and OWL detectors. Magnetic fields can lead to clustering of the predicted events with timescales spread from years to millions of years.

The second component stems from QNe with the compact remnant born in the propeller regime (pulsars with high magnetic field and small period undergoing a QN explosion immediately following a supernova) ejecting relativistic particles with $\Gamma > 1000$. The particles with $\Gamma_{\text{prop}} > \Gamma_{\text{QN}}$ will eventually interact with the QN shock, as to reach roughly $10^{16}$ eV (protons) and $10^{18}$ eV (iron) in one shock crossing. Given the energy of the QN shock, only a tiny amount of the propelled wind particles can be accelerated, i.e., a maximum of $10^{50}$ particles around the energy at the knee in Figure 1 (see eq. [14]). The propelled particles that were not accelerated by the QN shock are injected into the galactic space and will eventually become accelerated by the supernova shocks in the galaxy. QNe as possible UHECR sources seem to account for the observed extragalactic flux and can contribute partially (a few percent) to the galactic cosmic rays. We conclude by stating that despite the fact that there is no guarantee that QNe occur in nature, the model possesses features that can be tested in future cosmic-ray detectors.

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