ON THE ORIGIN AND SURVIVAL OF ULTRA-HIGH-ENERGY COSMIC-RAY NUCLEI IN GAMMA-RAY BURSTS AND HYPERNOVAE

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ABSTRACT

The chemical composition of the ultra-high-energy (UHE) cosmic rays serves as an important clue to their origin. Recent measurements of the elongation rates by the Pierre Auger Observatory hint at the possible presence of heavy or intermediate-mass nuclei in the UHE cosmic rays. Gamma-ray bursts (GRBs) and hypernovae have been suggested as possible sources of the UHE cosmic rays. Here we derive constraints on the physical conditions under which UHE heavy nuclei, if they are accelerated in these sources, can survive in their intense photon fields. We find that in the GRB external shock and hypernova scenarios, UHE nuclei can easily survive photodisintegration. In the GRB internal shock scenario, UHE nuclei can also survive, provided the dissipation radius and/or the bulk Lorentz factor of the relativistic outflow are relatively large, or if the low-energy self-absorption break in the photon spectrum of the prompt emission occurs above several keV. In internal shocks and in the other scenarios, intermediate-mass UHE nuclei have a higher probability of survival against photodisintegration than UHE heavy nuclei such as Fe.

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1. INTRODUCTION

Ultra-high-energy (UHE) cosmic rays above ~10^{18} eV are thought to be of extragalactic origin, since charged particles with such high energies cannot be confined by the magnetic field of our Galaxy. Galactic cosmic-ray accelerators, such as supernova remnants, are expected to reach their maximum energy below ~10^{18} eV. The transition between Galactic and extragalactic cosmic rays is therefore believed to happen either at the "second knee" around 10^{18} eV, where the chemical composition changes significantly as measured by HiRes, followed by a "dip" in the spectrum (Berezinsky et al. 2006), or at the "ankle" around 10^{19} eV, where the cosmic-ray spectrum becomes flatter.

At these ultrahigh energies, the chemical composition is a subject of debate. It has been speculated that these cosmic rays are made of pure protons, up to the highest energies. On the other hand, there are also reasons for favoring a cosmic-ray spectrum dominated by heavy or intermediate mass nuclei at the highest energies, since according to the Hills criterion (Hillas 1984), astrophysical sources are able to accelerate particles up to a maximum energy proportional to their nuclear charge. Recently, a mixed composition scenario has been invoked to fit the UHECR spectrum above ~10^{18.5} eV (Allard et al. 2005). By studying the details of the development of the resulting air showers, one can in principle infer the species of the primary UHECRs, since at a given energy, showers initiated by heavy nuclei develop higher up in the atmosphere than proton-induced showers. Both AGASA and HiRes data favor a dominance of light hadrons, consistent with being pure protons, in the composition of UHECRs above 10^{19} eV (Hughes et al. 2007), which is consistent with models assuming that UHECRs above 10^{18} eV are due to extragalactic protons. On the other hand, a recently report of an analysis of the muon component of UHE air showers in the Yakutsk data showed a heavy nuclei component in the UHECR spectrum (Glushkov et al. 2007).

One of the aims of Pierre Auger Observatory is to study the composition of UHECRs. It provides more precise measurements of the depth of the UHECR-initiated shower maximum X_{max} at extremely high energies (>10^{19} eV), with uncertainties that are about a factor 4 smaller than those of the best measurements from the HiRes experiment (Unger et al. 2007). The elongation rate data presented by the Pierre Auger Observatory team are satisfactorily represented by a fit containing a break point in the slope at 10^{18.35} eV (Unger et al. 2007). Due to the uncertainties in the hadronic interactions at the highest energies, the interpretation of these elongation rate depends on the hadronic physics models used in the analysis, and is therefore rather ambiguous at present. However, regardless of which interaction models are used, the data appear to require the presence of a substantial fraction of heavy or intermediate-mass nuclei in the range of Greisen-Zatsepin-Kuzmin (GZK) cutoff energies. A possible conflict in the chemical composition between the above Auger elongation results and recent results by the Pierre Auger collaboration (2007) on large-scale spatial correlations is an issue which remains to be resolved.

The proposed astrophysical source models of UHECRs include active galactic nucleus (AGN) jets (e.g., Biermann 1987; Berezinsky et al. 2006), intergalactic accretion shocks (e.g., Inoue et al. 2005), and gamma-ray bursts (e.g., Waxman 1995, 2004a; Vietri 1995; Wick et al. 2004; Dermer & Atoyan 2006; Murase et al. 2006). Recently, we have proposed that extragalactic hypernovae associated with subenergetic GRBs are also a probable source for UHECRs, whose energies are sufficient large to account for UHECRs from the second knee and above (Wang et al. 2007). These UHECR accelerators also accelerate electrons, which produce optical, X-ray, or gamma-ray photons through synchrotron and inverse Compton emission. Since a sufficiently high density of these low-energy photons can potentially lead to photodisintegration of the cosmic-ray heavy nuclei, a natural question is under what conditions these sources allow the survival of UHE
heavy nuclei from the acceleration site. In this paper we explore this problem for GRBs and hypernovae sources, since they have the most compact source sizes and might therefore pose the greatest hurdle for UHE heavy nuclei survival.

Long-duration GRBs are generally believed to result from the core collapse of massive stars, direct evidence for which comes from the spectroscopic identification of bright supernovae in association with these GRBs (see Woosley & Bloom 2006 for a review). The collapse leads to a black hole (or magnetar) plus accretion disk system. The annihilation of neutrinos or Poynting flows arising from the inner hot accretion disk produce a high-entropy fireball outflow, which expands and converts its internal energy into the kinetic energy of a small amount of baryon material, while propagating through the star along the rotation axis of the collapsing core. The outflow gets collimated, and finally two highly relativistic jets break out of the stellar envelope. The GRB scenario for UHECRs suggests that both internal shocks (Waxman & Atoyan 2006) and external shocks (Vietri 1995; Wick et al. 2004; Dermer & Atoyan 2006) in the jets are able to accelerate baryons to ultrahigh energies. Internal shocks develop due to internal collisions between the shells in the fireball jet ejecta, while the external shocks occur when the fireball ejecta interacts with the surrounding interstellar medium (see Mészáros 2006, Zhang 2007 for recent reviews). In the internal shock scenario, viewed as a source for the cosmic-ray nuclei, the jets themselves must contain heavy nuclei, whereas in external shocks the heavy nuclei could come from the swept-up circumstellar material that the jet is running into, presumably the stellar wind or interstellar medium. Since the two shocks have quite different dissipation radii and emission properties, the disintegration problem of the UHE cosmic-ray nuclei are different, as discussed in §§ 4 and 5, respectively.

Hypernovae are a peculiar type of supernovae with higher ejection velocities of the remnant shell and generally larger explosion energies than typical supernovae (Paczyński 1998). The prototype of hypernovae is SN 1998bw, a Type Ic supernova associated with an underenergetic GRB, GRB 980425 (Galama et al. 1998). Mildly relativistic ejecta components are inferred to be present in all three of the well-identified hypernova/underenergetic GRB systems so far, namely SN 1998bw/GRB 980425, SN 2003lw/GRB 031203, and SN 2006aj/GRB 060218. Based on this mildly relativistic ejecta component and the event rates of these objects, we have shown that hypernovae can accelerate particles to $\sim Z \times 10^{19}$ eV, and the energetics and occurrence rate is sufficient to account for the flux of UHECRs above the second knee ($\sim 6 \times 10^{17}$) eV, where Z is the nuclear charge of the accelerated particles (Wang et al. 2007). The particles are accelerated in the hypernova blast wave formed by the interaction between the hypernova ejecta and the surrounding stellar wind medium. In § 6 we study whether UHE heavy nuclei accelerated in such hypernovae can survive their photon environment.

The goal of this paper is to explore the possibility of acceleration of UHE nuclei in GRBs or hypernovae, the possible origins of these nuclei, and under what conditions they would survive their source environment; or conversely, under what conditions one could expect these sources to accelerate mainly protons to UHE energies.

2. ORIGIN OF HEAVY OR INTERMEDIATE-MASS NUCLEI IN THE SOURCES

It is at present unknown whether heavy or intermediate-mass nuclei are present in GRB jets. At the base of the outflow, the jets start out as a hot fireball, within a region of size $10^6$–$10^7$ cm and temperature $kT = 1$–$10$ MeV, where any heavy nuclei will be photodisintegrated due to the abundance of photons with energies comparable to the nuclear binding energy, $\sim 10$ MeV. The fireball is thus initially made up of free nucleons, $e^+\bar{e}$ pairs, trapped blackbody radiation and magnetic fields.

As the fireball expands and cools, the free nucleons in the jet may recombine into $\alpha$-particles (Beloborodov 2003), but they will not form heavy nuclei. However, as the jet burrows through the stellar core (from inside outward, Fe, O, and C cores in sequence), heavy nuclei from the stellar surroundings could be entrained into the jet. According to numerical simulations of jet propagation (Zhang et al. 2003), Kelvin-Helmholtz instabilities and/or oblique shocks that develop lead to the mixing of surrounding material into the jet, while the jet is advancing with a subrelativistic velocity. Since the temperature of the thermal photons trapped in the jet decreases significantly as the jet expands, with $T(r) \sim r^{-1}$, these thermal photons are no longer able to disintegrate the entrained nuclei when the jet has reached the radius of the stellar Fe core at $\sim 10^9$ cm. To avoid spallation of the nuclei due to nucleon-nucleon collisions, the relative velocity between the jet and the surrounding core material should be below a critical value $\beta_{sp} \approx 0.14$ (in units of the speed of light), at which the relative kinetic energy equals the nuclear binding energy, $\sim 10$ MeV. The calculations show that the jet head moves with a velocity of about $10^9$ cm s$^{-1}$ inside the Fe core at $\sim 10^9$ cm, and then the velocity increases as $\sim r^{1/2}$ (Mészáros & Rees 2001). Thus, the Fe and O nuclei entrained from the surroundings can survive both photodisintegration and nuclear spallation.

After the jet head breaks out of the star, in its subsequent stages the jet continues to escape through the evacuated funnel cavity through the star, so it accelerates to a relativistic velocity at radii still inside the star. Thereafter, the Kelvin-Helmholtz instability which causes the mixing of nuclei will be suppressed due to the relativistic relative motion (Zhang et al. 2003). Moreover, since the relative velocity between the jet and the surrounding star exceeds the critical velocity $\beta_{sp}$, spallation effects will need to be taken into account. One finds that the time for spallation $t_{sp} = 1/(\sigma_{sp} n_N c)$ is much shorter than the dynamic time $r/(\Gamma c)$ when the relativistic jet is inside the progenitor star, where $n_N = L/(4\pi r^2 \Gamma^2 m_N c^3)$ is the nucleon density of the jet, and $\Gamma$ is the Lorentz factor of jet. This means that even if nuclei get entrained, they will be spalled into lighter nuclei. Since the break-out time for initial jet after its inception is about $t_{b} = 10$ s, long GRBs with rest-frame durations $t_{\gamma} \sim 2$–$10$ s will have spent most of their life inside the star, while the jet head was moving subrelativistically, so most of the jet can become Fe-enriched before the internal shocks occur. For longer duration jets, with $t_{\gamma} \gtrsim 10$ s, only the first $\sim 10$ s portion of the burst may be Fe-rich, while the rest of the outburst may consist of pure protons.

The above discussion is relevant to the models that invoke the composition of the GRB jet as UHECR nuclei, which are accelerated in internal shocks that are widely assumed to be responsible for the prompt $\gamma$-ray emission. For GRB external shocks and hypernova remnant blast wave shocks, the nuclei could be due to the material swept-up by the shock front. The progenitors of long-duration GRBs and hypernovae are thought to be Wolf-Rayet stars, as the spectral type of the discovered supernovae in these events is typically Ic. These stars are stripped of their hydrogen envelope and sometimes even the helium envelope. The heavy or intermediate-mass UHE nuclei may originate from the stellar wind of the Wolf-Rayet star. In WC type Wolf-Rayet stars, the C abundance is $X_C = 20\%$–$55\%$ (by mass) and the O abundance is $X_O = 5\%$–$10\%$ (Crowther 2007). In WO type Wolf-Rayet stars,
O and C abundances are even higher, with \(X_O = 15\%–25\%\) and \(X_C = 40\%–55\%\) (e.g., Kingsburgh et al. 1995). The abundance of heavy elements in stellar wind of these Wolf-Rayet stars is clearly much higher than the solar value.

### 3. DISINTEGRATION OF UHE NUCLEI

The most relevant processes that may prohibit acceleration of heavy nuclei to UHE at the astrophysical sources are photodisintegration, photopion production, and nuclear spallation. The first two processes are due to interactions of UHE nuclei with surrounding photons, and the third is due to interactions with other nuclei or nucleon. Both the photodisintegration and photopion processes are dominated by resonance production of either an excited state of the nuclei or a delta and subsequent de-excitation or decay. In case of photodisintegration, this dominant channel is called a giant dipole resonance (GDR) in the energy range \(\sim 10–30\) MeV, with a threshold energy of \(\sim 10\) MeV per nucleon. In case of photopion the threshold energy for pion production is \(\simeq 150\) MeV.

For an UHE nucleus with Lorentz factor \(\gamma_A\) propagating through an isotropic photon background with differential number density \(n(\varepsilon)\) at energy \(\varepsilon\), the photodisintegration or photopion rate is given by (Stecker 1968)

\[
\tau_{\text{dis},p}^{-1} = \frac{\gamma_A}{e} \int_{\varepsilon_0}^{\varepsilon_0 + \Delta \varepsilon} \frac{n(\varepsilon)}{\varepsilon^2} d\varepsilon, \tag{1}
\]

where \(\varepsilon' = \varepsilon - \varepsilon_0\) is the energy in the rest frame of the nucleus. The threshold energy \(\varepsilon_0\) for photodisintegration and photopion is \(10\) MeV. The threshold energy \(\varepsilon_0\) is \(15\) MeV for photodisintegration and \(145\) MeV for photopion. The cross sections for photodisintegration in the energy range \(\varepsilon_0 < \varepsilon' \leq 30\) MeV with loss of one nucleon can be approximately described in a Lorentzian form (Puget et al. 1976; Anchordoqui et al. 2007) as

\[
\sigma(\varepsilon') = \frac{\sigma_0 \Delta \varepsilon^2 \Delta_{\text{GDR}}^2}{(\varepsilon_0^2 + \varepsilon'^2 + \Delta_{\text{GDR}}^2)^2}, \tag{2}
\]

where \(\Delta_{\text{GDR}}\) and \(\sigma_0\) are the width and maximum value of the cross section, and \(\varepsilon_0\) is the energy at which the cross section peaks. Fitted numerical values are \(\sigma_0 \simeq 1.45 \times 10^{-27}\) cm², \(\Delta_{\text{GDR}} = 8\) MeV, and \(\varepsilon_0 = 4.26 \times 10^{-21}\) MeV for \(A > 4\) (Karakuła & Tkaczyk 1993). Above \(\varepsilon' > 30\) MeV and below the pion production threshold energy, photodisintegration may result in multiple nucleon emission, although with a much lower cross section. Puget et al. (1976) suggested a parameterization such that the cross section \(\sigma(\varepsilon')\) integrated in the range \(\varepsilon_0 < \varepsilon' < 30\) MeV and in the range \(30\) MeV \(< \varepsilon' \leq 150\) MeV is equal. We assume a flat cross section in the range \(30\) MeV \(< \varepsilon' \leq 150\) MeV of \(6.6 \times 10^{-27}\) cm² and \(1.7 \times 10^{-27}\) cm², respectively, for iron and oxygen nuclei, satisfying the above condition. This more accurate cross section affects the photodisintegration rate when the photon spectrum is very hard, as we will see later. For soft photon spectra equation (2) is adequate, and often a delta function approximation \(\sigma(\varepsilon') \approx \sigma_0 \Delta_{\text{GDR}} (\gamma_A^2 - 1)\) can provide an order-of-magnitude estimate.

The photopion cross section formula for delta resonance production is well known and is parameterized by Mücke et al. (2000) as

\[
\sigma(\varepsilon') = \frac{\sigma_{0,\gamma} s^2 \Delta_\gamma^2}{\varepsilon' \left(\frac{m_p^2 c^4}{s} - s + s \Delta_\gamma^2\right)}, \tag{3}
\]

where \(s = m_p^2 c^4 + 2\varepsilon' m_p c^2\) is the center-of-mass energy, \(\sigma_{0,\gamma} = 3.11 \times 10^{-29}\) cm², and the peak cross section is \(\sim 4.12 \times 10^{-28}\) cm² at \(\varepsilon' \simeq 0.3\) GeV. The width of the resonance is \(\Delta_\gamma \simeq 0.11\) GeV. The photonuclear cross section is \(\sigma(\varepsilon' \simeq 0.11\) GeV.

The cross section for spallation for a nucleus with atomic number \(A\), \(n_N\) is the number density of nucleis in the sources, and \(\beta\) is the velocity of the UHE nucleons. As we show below, even for internal shocks, which have the highest nucleon density, this spallation effect is less important than photodisintegration.

Below we discuss the photodisintegration process in all three possible UHE nuclei sources, i.e., GRB internal shocks, GRB external shocks, and hypernova remnant blast waves.

### 4. GRB INTERNAL SHOCKS

Due to the variable injection at the base, internal collisions occur within the unsteady plasma jet, which develop into internal shocks. The prompt gamma-ray burst emission is supposed to result from the nonthermal emission of electrons accelerated in these shocks. The two main processes through which the heavy nuclei could be disintegrated in internal shocks are photodisintegration or photopion interaction by X-ray photons in prompt emission, and spallation due to collisions with other nucleons in the relativistic jet.

#### 4.1. Nucleus Photodisintegration due to Prompt X-Rays

The GRB prompt photon spectrum is well fitted in the BATSE range (10 keV to 3 MeV) by a combination of two power laws, \(n(\varepsilon) \propto \varepsilon^{-\beta}\), with different values of \(\beta\) at low and high energy (Band et al. 1993). The break energy \(\varepsilon_{\text{b,obs}}\) in the observer frame is typically \(\varepsilon_{\text{b,obs}} = \Gamma \varepsilon_{\text{b}} \sim 1\) MeV, with \(\beta \simeq 1\) at energies below the break and \(\beta \approx 2\) above the break, where \(\Gamma\) is the bulk Lorentz factor of the relativistic flow that develops internal shocks, and \(\varepsilon_{\text{b}}\) is the break photon energy in the flow rest frame. At low energy, in the framework of the internal shock model, there could be another spectral break, i.e., the synchrotron self-absorption break, \(\varepsilon_{\text{ssa}}\).

Thus the photon spectrum, in the comoving frame, is

\[
n(\varepsilon) = \begin{cases} 
n_b(\varepsilon/\varepsilon_{\text{b}})^{-2}, & \varepsilon > \varepsilon_{\text{b}}, \\
_b(\varepsilon/\varepsilon_{\text{b}})^{-1}, & \varepsilon_{\text{ssa}} < \varepsilon < \varepsilon_{\text{b}}, \\
_b(\varepsilon_{\text{ssa}}/\varepsilon_{\text{b}})^{-1}(\varepsilon/\varepsilon_{\text{ssa}}), & \varepsilon < \varepsilon_{\text{ssa}}, \end{cases} \tag{5}
\]

where \(n_b \equiv n(\varepsilon_{\text{b}})\) is the photon number density at the break energy in the comoving frame of the wind.

Using this photon spectrum, the inner integral in equation (1) gives

\[
\int_{\varepsilon/2\gamma}^{\infty} n(\varepsilon) \varepsilon^{-2} d\varepsilon = \frac{n_b}{(1 + \beta)\varepsilon_{\text{b}}} \left(\frac{\varepsilon}{2\gamma \varepsilon_{\text{b}}}\right)^{-(1 + \beta)}.
\]

The energy density of photons in the wind rest frame in the energy range of BATSE is \(U_\gamma = n_b \varepsilon_{\text{b}}^2 (1 + \ln(\varepsilon_{\text{b}}/\varepsilon_{\text{b}})) \approx 2n_b \varepsilon_{\text{b}}^2\), where \(\varepsilon_{\text{b}} = 3\) MeV is the upper energy range of BATSE. This energy density is related to the observed photon luminosity of GRBs,
whose typical value is \( L_\gamma \sim 10^{51} \text{ erg s}^{-1} \), by \( L_\gamma = 4\pi R_m^2 \Gamma^2 cU_\gamma \), where \( R_m \) is the radius of the emitting region of these photons.

The first integral of equation (1) can be done numerically for the exact cross sections in equations (2) and (3). For the sake of an analytic treatment, we approximate the cross section as being mainly contributed by the resonance peak and find that

\[
\tau_{\text{dis}}^{-1} = \begin{cases} 
\frac{U_{\gamma}}{2\varepsilon_b} \sigma_{\text{GDR}} \frac{\Gamma_5}{\varepsilon_b} \frac{\Gamma_2}{\varepsilon_0} \frac{2\varepsilon_b}{\varepsilon_0} & 2\varepsilon_b \leq \varepsilon_0, \\
\frac{U_{\gamma}}{2\varepsilon_b} \sigma_{\text{GDR}} \frac{\Gamma_5}{\varepsilon_b} \frac{\Gamma_2}{\varepsilon_0} & 2\varepsilon_{\text{ssa}} \leq \varepsilon_0 \leq 2\varepsilon_b, \\
\frac{U_{\gamma}}{2\varepsilon_b} \sigma_{\text{GDR}} \frac{\Gamma_5}{\varepsilon_b} \frac{\Gamma_2}{\varepsilon_0} \kappa^{-2} (2 \ln \kappa + 1) & 2\varepsilon_{\text{ssa}} \geq \varepsilon_0,
\end{cases}
\]

where \( \kappa \equiv 2\varepsilon_{\text{ssa}}/\varepsilon_0 \). The observer-frame energy of the nucleus that interacts preferentially with photons with energy \( \varepsilon_b \) and \( \varepsilon_{\text{ssa}} \) are

\[
E_b = 10^{17} \left( A/56 \right)^\frac{1}{2} (\varepsilon_b/1 \text{ MeV})^{-1} \text{ eV},
\]

\[
E_{\text{ssa}} = 10^{20} \left( A/56 \right)^\frac{1}{2}(\varepsilon_{\text{ssa}}/1 \text{ keV})^{-1} \text{ eV},
\]

where \( \Gamma \sim 300 \Gamma_5 \) is a typical GRB bulk Lorentz factor.

Now we can obtain the time \( t_{\text{dis}} \) of an nucleus of energy \( E \) in the observer frame: for \( E_b < E < E_{\text{ssa}} \),

\[
t_{\text{dis}} = 0.035 \left( \frac{A}{56} \right)^{-\frac{1}{21}} \left( L_{\gamma,5} R_{\text{in},13}^2 \frac{\varepsilon_b}{1 \text{ MeV}} \right) \text{ s},
\]

and for \( E > E_{\text{ssa}} \),

\[
t_{\text{dis}} \approx 0.035 \left( \frac{A}{56} \right)^{-\frac{1}{21}} L_{\gamma,5} R_{\text{in},13}^2 \frac{\varepsilon_{\text{obs}}}{1 \text{ MeV}} \left( \frac{E}{E_{\text{ssa}}} \right)^{\frac{2}{5}} \text{ s}.
\]

The dynamic time of internal shock is

\[
t_{\text{dyn}} = R_{\text{in}}/c\Gamma = 10 R_{\text{in},13} \Gamma_5^{-\frac{1}{2}} \text{ s}.
\]

By comparing equations (8) or (9) with equation (10), we can see that the heavy nuclei can survive when either

\[
R_{\text{in},13} \Gamma_5^2 \gg 30 L_{\gamma,51} \left( \frac{\varepsilon_{\text{obs}}}{1 \text{ MeV}} \right)^{-1} \left( \frac{A}{56} \right)^{\frac{1}{21}}
\]

corresponding to \( t_{\text{dis}} \gtrsim t_{\text{dyn}} \), or

\[
\varepsilon_{\text{ssa,obs}} \gtrsim 5 R_{\text{in},13}^{\frac{1}{2}} \Gamma_5 E_{\text{20}} \left( \frac{A}{56} \right)^{\frac{1}{21}} L_{\gamma,51}^{-\frac{1}{2}} \left( \frac{\varepsilon_{\text{obs}}}{1 \text{ MeV}} \right)^{-\frac{1}{2}} \text{ keV},
\]

in the case of \( t_{\text{dis}} \lesssim t_{\text{dyn}} \).

For comparison, the acceleration time of UHE nuclei with nuclear charge \( Z \) in the internal shock is

\[
t_{\text{acc}} = \frac{\alpha E}{ZVeBc} = 1.5 \times 10^{-2} \alpha E_{20} L_{k,5}^{-\frac{1}{2}} R_{\text{in},13} L_{\gamma,51}^{-\frac{1}{2}} \left( \frac{Z}{26} \right)^{-\frac{1}{2}} \text{ s},
\]

where \( \alpha \sim 1 \), describing the ratio between the acceleration time and Larmor time, and

\[
B = \left( \frac{8\pi e\mu_L}{4\pi R^2 \Gamma e} \right)^{\frac{1}{2}} = 8 \times 10^4 \left( \frac{1}{L_{k,5}} \right)^{\frac{1}{2}} R_{\text{in},13} \Gamma_5^{-\frac{1}{2}} \text{ G}
\]

is the comoving frame magnetic field in internal shocks. The synchrotron loss time for UHE nuclei is

\[
t_{\text{syn}} = \frac{6\pi m_e^4 c^3 \Gamma}{\sigma_T m_e^2 EB^2} \left( \frac{A}{Z} \right)^4 \sim 30 \left( \frac{\varepsilon_{\text{obs}}}{1 \text{ keV}} \right)^{-\frac{1}{2}} R_{\text{in},13}^{\frac{1}{2}} \Gamma_5^{\frac{3}{2}} E_{20}^{-\frac{1}{2}} \text{ s},
\]

which is a factor of \( (A/Z)^4 \) \sim 16 longer than that of protons.

By equating \( t_{\text{acc}} \) with \( t_{\text{syn}} \) or \( t_{\text{acc}} \), we get the maximum energies of accelerated nuclei for these two cases, respectively:

\[
\varepsilon_{\text{max}} = 6 \times 10^{21} \varepsilon_{\text{obs}}^{-\frac{1}{2}} R_{\text{in},13}^{\frac{1}{2}} L_{k,52}^{-\frac{1}{2}} \left( \frac{Z}{26} \right)^{\frac{1}{2}} \text{ eV},
\]

\[
\varepsilon_{\text{max}} = 4 \times 10^{21} \varepsilon_{\text{obs}}^{-\frac{1}{2}} R_{\text{in},13}^{\frac{1}{2}} L_{k,52}^{-\frac{1}{2}} \left( \frac{Z}{26} \right)^{\frac{1}{2}} \text{ eV}.
\]

We compare these timescales in Figure 1 for two different sets of parameters, i.e., \( R = 10^{13} \text{ cm} \), \( \Gamma = 10^{2.5} \) (top panel) and \( R = 10^{14} \text{ cm} \), \( \Gamma = 10^3 \) (bottom panel). We can see that (1) the iron
nucleus can be accelerated to energies $\gtrsim 10^{20}$ eV, since in both cases, $t_{\text{acc}} < t_{\text{dyn}} < t_{\text{syn}}$ for $E = 10^{20}$ eV, and (2) when the internal shock radius and/or the bulk Lorentz factor are relatively large, the photodisintegration time of iron nuclei is longer than the dynamic time, as seen in the bottom panel of Figure 1. Note that the photodisintegration time and the photopion interaction times in these figures are calculated numerically according to equation (1) with exact cross sections.

For smaller internal shock radii (Fig. 1, top panel), the photodisintegration effect becomes important if the photon spectrum $n(\varepsilon) \sim \varepsilon^{-1}$ extends to $\lesssim 1$ keV. This is consistent with the result in Anchordoqui et al. (2008). However, if there is a self-absorption break at several keV in the photon spectrum that leads to a drop in the number density of the photons with which the UHE nucleus mainly interacts, the photodisintegration effect will be suppressed, as can be seen from the dashed lines, where a synchrotron self-absorption break at 5 keV in the photon spectrum has been assumed. This is possible for a small internal shock radius, since a smaller internal shock radius might lead to higher synchrotron self-absorption break $\varepsilon_{\text{ssa}}$. The value of $\varepsilon_{\text{ssa}}$ also depends on the dissipation model, as well as the radiation mechanism for prompt emission (e.g., Panaitescu & Mészáros 2000; Rees & Mészáros 2005; Pe’er et al. 2005; Pe’er & Zhang 2006), which is, however, largely uncertain at present. For the simple optically thin internal shock model, $\varepsilon_{\text{ssa}}$ is estimated to be

$$\varepsilon_{\text{ssa, obs}} \sim 1 L_{k, 52}^{1/3} \varepsilon_{\text{e}^{-1}}^{-1/3} R_{\text{i}, 13}^{-2/3} \Gamma_{2.5}^{-1/3} \text{keV}. \quad (16)$$

For UHE intermediate-mass nuclei, such as O nuclei, the constraints for survival, equations (11) and (12), are much looser. Thus, for small internal dissipation or shock radii, intermediate-mass nuclei are much more likely to survive photodisintegration. This can be also seen in Figure 2, where all the timescales are calculated for O nuclei with $Z = 8$ and $A = 16$.

4.2. Nuclear Spallation due to Collisions

An UHE nucleus can also spallate due to collision with other nucleons in internal shocks. The nucleon density of a relativistic flow with kinetic energy luminosity $L_k = 10^{52}$ erg s$^{-1}$ at internal shock radius $R_{\text{in}}$ is

$$n_N = \frac{L_k}{4\pi R_{\text{in}}^2 \Gamma^2 m_N \varepsilon^{3}} = 1.8 \times 10^{12} L_{k, 52} R_{\text{in}, 13}^{-2} \Gamma_{2.5}^{-2} \text{ cm}^{-3}.$$  

So the spallation time is

$$t_{\text{sp}} = \frac{1}{n_N \sigma_{\text{sp}} \varepsilon} = 25 L_{k, 52} R_{\text{in}, 13} \Gamma_{2.5}^{2} \left( \frac{A}{56} \right)^{-2/3} \text{ s}. \quad (17)$$

Comparing this with the dynamic time, we find that the heavy nucleus can survive if

$$R_{\text{in}, 13} \Gamma_{2.5}^{3} \gtrsim 0.04 L_{k, 52} \left( \frac{A}{56} \right)^{2/3} . \quad (18)$$

This condition is much easier to satisfy than the more stringent photodisintegration constraint of equation (11).

5. GRB EXTERNAL SHOCKS

As the relativistic GRB jets expand and sweep up the surrounding material, external shocks develop. The accelerated electrons in the external shock produce the observed afterglow emission. It has also been suggested that the external shock may be able to account for UHECRs. The early deceleration phase, when the jet converts about half of its energy to the swept-up material, is the most effective phase for the external shock to accelerate particles, since the shocked material has the largest Lorentz factor at that time. The early external shock produces the early X-ray afterglow, whose average luminosity observed by Swift is about $L_X \approx 10^{46}$ erg s$^{-1}$ during the first $\sim 10$–100 s (Nousek et al. 2006). In this external shock scenario, the accelerated UHE nuclei with energy $E \gtrsim 10^{19}$ eV preferentially interact with these X-ray photons, if the Lorentz factor is typically $10^2 \lesssim \Gamma \lesssim 10^3$.

Taking the early X-ray afterglow spectrum in the fast-cooling regime (e.g., Mészáros 2006) with $F_{\nu} \sim \nu^{-1}$, corresponding to photon spectral index $\beta = 2$, we can obtain the photodisintegration rate of a nucleus moving with Lorentz factor $\gamma_A$, i.e.

$$t_{\text{dis}}^{-1} = \frac{4}{3} \sigma_0 \frac{\Delta_{\text{GDR}}}{\varepsilon_0} \frac{\gamma_A L_X}{\kappa \varepsilon_0}, \quad (19)$$

where $L_X$ is the comoving-frame energy density of X-ray afterglow photons and $\kappa = \ln(\epsilon_{\text{XRT,M}}/\epsilon_{\text{XRT,m}}) \approx 3$, where $\epsilon_{\text{XRT,M}}$ and $\epsilon_{\text{XRT,m}}$ being the upper and lower end of Swift XRT energy threshold. The energy density is related to the luminosity by $L_X = 4\pi R_{\text{ex}}^2 \Gamma^2 c \epsilon_U$, where $R_{\text{ex}}$ is the radius of the external shock at the end of the free expansion phase of the ejecta. Now we can obtain the photodisintegration energy loss time of a cosmic-ray nucleus of energy $E$ propagating in the early afterglow photons,

$$t_{\text{dis}} = 3 \times 10^6 R_{\text{ex}, 17} \Gamma_{2.5}^{-1} \epsilon_{\text{X},48} E_{20}^{-1} \left( \frac{A}{56} \right)^{-0.21} \text{ s}. \quad (20)$$
Comparing this with the dynamic time in the comoving frame of external shocks $\tau_{\text{dyn}} = R_{\text{ex}}/(2\epsilon) = 10^8 R_{\text{ex},17} \Gamma^{-1}_{2.5} s$, we can see that the heavy nucleus can survive when

$$R_{\text{ex},17} \Gamma^{-4}_{2.5} \gtrsim 3 \times 10^{-3} L_{X,48} E_{20} \left( \frac{A}{56} \right)^{0.21}.$$ (21)

For a constant density medium environment, this gives an upper limit on the number density of the medium,

$$n \lesssim 2.5 \times 10^6 E_{k,53} \Gamma^{10}_{2.5} L_{X,48}^{-1} E_{20}^{-1} \left( \frac{A}{56} \right)^{-0.63} \text{ cm}^{-3},$$ (22)

while for a stellar wind environment with a density profile

$$\rho = \left( \frac{\dot{M}}{4\pi v_w} \right) r^{-2} = 5 \times 10^{11} A_r r^{-2} \text{ g cm}^{-1},$$

this gives constrains on the mass-loss rate parameter

$$A_r \lesssim 1.3 E_{k,53} \Gamma^{-2}_{2.5} L_{X,48}^{-1} E_{20}^{-1} \left( \frac{A}{56} \right)^{-0.21},$$ (23)

where $A_r = 1$ corresponds to $\dot{M} = 10^{-5} M_\odot \text{ yr}^{-1}$, and $v_w = 1000 \text{ km s}^{-1}$.

After the initial free expansion phase, the external shock starts to decelerate as more and more material is swept up. The shock radius evolves with time as $R_{\text{ex}}(t) = 4 \times 10^{17} E_{k,53}^{1/4} A_r^{1/4} t_d^{1/4} \text{ cm}$ for constant-density medium and $R_{\text{ex}}(t) = 2 \times 10^{18} E_{k,53}^{1/2} A_r^{-1/2} t_d^{1/2} \text{ cm}$ for a wind medium, where $t_d$ is time in units of days. Assuming that the magnetic field energy density acquires a fraction $\epsilon_B = 0.1$ of the internal energy, the magnetic field evolves as $B = 1.5 \epsilon_B^{-1/2} E_{k,53}^{1/8} A_r^{1/4} t_d^{1/4} \text{ G}$ for a constant-density medium and $B = 0.07 \epsilon_B^{-1/2} E_{k,53}^{1/2} A_r^{-1/2} t_d^{1/2} \text{ G}$ for a wind medium. Since the synchrotron loss is typically unimportant for UHE nuclei in external shocks, the maximum energy of particles accelerated in external shocks is obtained by equating $\tau_{\text{sec}}$ with $\tau_{\text{dyn}}$, which gives

$$\epsilon_{\text{max}} = Z e B R_{\text{ex}}(t) = \begin{cases} 3 \times 10^{21} \frac{Z}{26} \epsilon_B^{-1/2} E_{k,53}^{3/8} A_r^{1/4} t_d^{1/4} \text{ eV ISM,} \\ 10^{-31} \frac{Z}{26} \epsilon_B^{-1/2} E_{k,53}^{1/4} A_r^{1/4} t_d^{1/4} \text{ eV wind.} \end{cases}$$ (24)

This shows that the maximum energy of accelerated particles by external shock decreases rather slowly with time, so that even weeks to months after the burst, heavy nuclei can be still accelerated to UHE energies by afterglow shocks, even if the flux may decrease with time.

The optical depth $\tau$ for photodisintegration of these nuclei evolves with time as

$$\tau = \frac{\tau_{\text{dyn}}}{\tau_{\text{dis}}} \sim L_X \Gamma^{-4} R_{\text{ex},t}^{-1} \sim \begin{cases} t_d^{1/4} \text{ ISM,} \\ t_d^{-1/2} \text{ wind,} \end{cases}$$ (25)

where an X-ray afterglow luminosity $L_X \sim t^{-1}$ has been assumed. The slowly increasing or decreasing optical depth indicates that the accelerated heavy nuclei by GRB external shocks can survive for a relatively long time. We plot in Figure 3 the evolution of photodisintegration optical depth and the maximum energy of accelerated particles for iron nuclei (top panel) and oxygen nuclei (bottom panel), respectively.

![Figure 3](image-url)

**Fig. 3.**—External shocks scenario. Top: Dashed lines show the evolution of the maximum energy of an accelerated iron nucleus for an external shock moving into a constant-density medium (thick dashed line) and into a wind medium (thin dashed line), while the solid lines show the evolution of the photodisintegration optical depth of an iron nucleus of energy $E = 10^{20}$ eV in external shocks moving into the constant-density medium (thick solid line) and the wind medium (thin solid line). The parameters used for the constant-density medium case and the wind medium case are, respectively, ($n = 1 \text{ cm}^{-3}, E = 10^{33} \text{ erg}, \Gamma = 10^{3}, \epsilon_B = 0.1$), and ($n = 0.1, E = 10^{53} \text{ erg}, \Gamma = 10^{2.5}, \epsilon_B = 0.1$). Bottom: Same as the top panel, but for UHE oxygen nuclei.

6. HYPERSONA REMNANT BLAST WAVE

The observations of the radio afterglow of the hypernova SN 1998bw showed that about $10^{50}$ erg of kinetic energy were released in the form of a mildly relativistic ejecta (Kulkarni et al. 1998; Li & Chevalier 1999). The interpretation of the X-ray afterglow also favors a mildly relativistic ejecta component (Waxman 2004b). A recently detected strong thermal X-ray emission component in another subenergetic burst (GRB 060218), associated with SN 2006aj, may also be associated with a mildly relativistic supernova shock breakout, in which the mildly relativistic supernova ejecta has an energy $\sim 10^{49}$ erg (Campana et al. 2006). Due to the large supernova explosion energy and the much lower than typical GRB energy, attempts have been made to ascribe the prompt gamma-ray emission to the shock from the mildly relativistic ejecta as it breaks out through the hypernova progenitor’s outer envelope (Woosley et al. 1999; Tan et al. 2001), although a generally accepted conclusion has not yet been reached. We have used the term semirelativistic hypernovae to denote such supernovae exhibiting a mildly relativistic ejecta component, seen in association with GRBs. Such high-velocity ejecta are a key ingredient of the hypernova model for UHECRs, since the low-velocity bulk ejecta is not able to accelerate particles to such high energies. Based on the shock breakout scenario, we have suggested that there might be a continuous distribution of ejecta energy in velocity, i.e., $E_k(\Gamma \beta) \approx 3 \times 10^{52}\Gamma \beta/0.1 \alpha^{-\alpha}$ with $\alpha \approx 2$, $\Gamma \approx 5$. So, in this energy range, the bulk of the kinetic energies contained in the mildly relativistic supernova shock breakout ejecta is at least $10^{49}$ erg.
where \( \beta \) is the ejecta velocity in units of speed of light (Wang et al. 2007).

As the hypernova ejecta expand, they transfer their energy to the swept-up stellar wind, and external shocks develop, which accelerate wind particles to ultrahigh energies. Since the ejecta has a velocity distribution profile, the leading edge higher velocity ejecta decelerate the earliest, and then the lower velocity ejecta decelerate progressively. The maximum energy of accelerated particles is related to the ejecta velocity \( \Gamma \beta \) by

\[
\varepsilon_{\text{max}} \approx Z e BR \beta \Gamma \\
= 1.3 \times 10^{20} \frac{Z}{26} e^{1/2} \left( \frac{\Gamma \beta}{0.5} \right)^2 A_*^{1/2} \text{ eV.} \tag{26}
\]

Although higher velocity ejecta can accelerate particles to higher energies, the kinetic energy in ejecta of \( \Gamma \beta > 0.5 \) are too low to account for the UHECR flux (Wang et al. 2007), due to the steep distribution of energy \( E_\gamma \propto (\Gamma \beta)^{-2} \). For ejecta with a velocity \( \Gamma \beta \), the free expansion phase before deceleration sets in lasts for a time

\[ t_f = 1300(\Gamma \beta/0.5)^{-3} A_*^{-1} \text{ days,} \tag{27} \]

and the radius of the ejecta at this time is

\[ r_f = 1.7 \times 10^{18}(\Gamma \beta/0.5)^{-4} A_*^{-1} \text{ cm.} \tag{28} \]

There are two photon sources that could cause photodisintegration of heavy nuclei: one is provided by hypernova thermal photons from radioactive elements of the hypernova ejecta, and another is the synchrotron photons from the hypernova remnant blast wave. For the first mechanism, we use the luminosity of SN 1998bw as a representative. At time \( t \approx 1000 \) days after the burst, the optical luminosity of SN 1998bw drops to the level of about \( L_{\text{HE}} \sim 10^{40} \text{ erg s}^{-1} \) (Sollerman et al. 2002). A nucleus of energy \( E = 10^{20} \) eV interacts with target photons with energy \( \varepsilon_\gamma \gtrsim 0.01(A/56)E_{20} \) eV. A rough estimate of the optical depth of photodisintegration of heavy nuclei due to hypernova thermal photons is

\[
\tau \approx \sigma_0 \frac{L_{\text{HE}}}{4\pi f r_f} \frac{r_f}{\eta} = 3 \times 10^{-5} L_{\text{HE,30}} r_f^{-1} \frac{\varepsilon_{\text{HE}}^{1/2}}{1 \text{ eV}} \tag{29}
\]

where \( \eta \approx 4 \) is the compression ratio of the hypernova external shock and \( \varepsilon_{\text{HE}} \approx 1 \) eV is the characteristic energy of hypernova thermal photons. We can see that thermal photons are so sparse that they have a negligible effect on the photodisintegration of UHE nuclei.

The synchrotron emission from the stellar wind shocked by the hypernova remnant could be brighter than the hypernova emission itself at the late stages of the hypernova, e.g., \( t_f \sim 1300 \) days after the burst in our case, since the hypernova luminosity drops rather quickly after the peak. We estimate the luminosity from the shocked wind at the time when the \( \Gamma \beta = 0.5 \) ejecta begins to decelerate. The total number of the swept-up wind particles is

\[
N_e = \left( \frac{M}{v_w} \right) r_f = 6.6 \times 10^{54} \frac{t_f}{1300 \text{ days}} A_*^{-4/5} \text{ cm}^{-3}.
\]

Assuming that the magnetic field energy density is amplified to a fraction \( \varepsilon_B \) of the shock internal energy, the magnetic field is

\[
B = 0.015(\Gamma \beta/0.5)^2 A_*^{3/2} \text{ G} \approx 0.015 \left( \frac{t_f}{1300 \text{ days}} \right)^{-1/2} A_*^{3/2} \text{ G.}
\]

Assuming a power-law energy distribution for accelerated electrons \( d\nu/d\gamma_e \propto \gamma_e^{-p} \) with \( p \sim 2 \), one obtains the characteristic frequency

\[
\nu_m = 7 \times 10^5 \varepsilon_{e,-1}^2 \left( \frac{t_f}{1300 \text{ days}} \right)^{-9/5} A_*^{3/10} \text{ Hz.}
\]

where \( \varepsilon_{e,-1} \) is the fraction of the shock energy that goes into electrons. High-energy electrons will cool in the magnetic field and cause a break in the electron energy distribution. The characteristic synchrotron frequency corresponding to this break is

\[
\nu_e = 4 \times 10^{13} \left( \frac{t_f}{1300 \text{ days}} \right) A_*^{-3/2} \text{ Hz.}
\]

The luminosity at the frequency \( \nu_m \) is

\[
L_{\nu_m} = N_{\nu,m} P(\gamma_{e,m}) \approx 4 \times 10^{36} \varepsilon_{e,-1}^2 \varepsilon_{B,-1} \left( \frac{t_f}{1300 \text{ days}} \right)^{-2} A_* \text{ erg s}^{-1}.
\]

For \( p \sim 2 \), the luminosity at frequency \( \nu \) is \( L_\nu = L_{\nu_m}(\nu/\nu_m)^{1/2} \) for \( \nu_m < \nu < \nu_e \), while for \( \nu > \nu_e \), \( L_\nu \approx L_{\nu_m}(\nu/\nu_m)^{1/2} \). The luminosity at energy \( \varepsilon_{\gamma} = 0.01(A/56)E_{20} \) eV, corresponding to the energy of the photons with which the UHE nuclei of energy \( E \) interact at the resonance peak, is

\[
L_{\text{syn}} = 10^{40} \varepsilon_{e,-1}^{3/4} \varepsilon_{B,-1}^{1/2} \left( \frac{t_f}{1300 \text{ days}} \right)^{-1/1.9} A_*^{0.71} \text{ erg s}^{-1}.
\]

So we see that for wind parameters \( A_* = 1 \), the synchrotron luminosity does exceed the hypernova luminosity at the time \( \sim 10^3 \) days.\(^6\) Once we know the synchrotron luminosity of a hypernova remnant and the photon spectral index \( \beta = 3/2 \), we can, using equation (1), get the optical depth of photodisintegration of heavy nuclei due to such synchrotron photons, i.e.,

\[
\tau \approx 8 \times 10^{-3} \varepsilon_{e,-1}^{3/4} \varepsilon_{B,-1}^{1/2} A_*^{1.35} E_{20}^{1/2} \left( \frac{A}{56} \right)^{0.71} \left( \frac{t_f}{1300 \text{ days}} \right)^{-1.9}.
\]

Using \( \tau \approx 1 \), we find that, after a time

\[ t_f \gtrsim 100 \varepsilon_{e,-1}^{0.53} \varepsilon_{B,-1}^{0.4} A_*^{0.7} E_{20}^{0.26} \left( \frac{A}{56} \right)^{0.37} \text{ days,} \tag{31} \]

the nucleus can survive in the hypernova synchrotron photon environment. Note that the free expansion time for ejecta with \( \Gamma \beta \lesssim 0.5 \) (eq. [27]) is longer than this time. The maximum energy of accelerated particles, equation (26), depends on the time \( t_f \) as

\[
\varepsilon_{\text{max}} \approx 1.3 \times 10^{20} \frac{Z}{26} \left( \frac{t_f}{1300 \text{ days}} \right)^{1/2} A_*^{-2/5} \text{ eV.} \tag{32}
\]

\(^6\) Unlike in SN 2006aj, the inferred stellar wind for SN 1998bw is much weaker, with \( A_* = 0.04-0.1 \), which may explain why the optical emission of SN 1998bw, detected at \( \sim 1000 \) days after the burst, is still contributed by the hypernova emission.
The evolution of the optical depth for photodisintegration and the maximum energy of accelerated nuclei are shown in Figure 4.

7. DISCUSSION AND CONCLUSIONS

The possible presence of heavy nuclei in UHECRs raises interesting questions: What is the origin of these nuclei? Can these nuclei survive in the sources where they get accelerated? In this paper, we have endeavored to address these questions for two proposed UHECR sources discussed in the literature, namely, GRBs and hypernovae. For GRBs, both internal and external shocks have been suggested to accelerate particles to UHE energies, and in this paper we have considered the role of both of these in the context of UHECR heavy nuclei. We have sketched out some possible mechanisms for the presence of heavy nuclei in these scenarios. For GRB internal shocks, we suggest that the nuclei are entrained from the progenitor stellar core, e.g., the Fe core, the O core, etc., during the stage when the jets arising from accretion in the innermost collapsing core are making their way out through the star. Since instabilities favoring entrainment are predominant mainly during the initial stellar crossing phase of the jet, a jet composition with a substantial fraction of heavy nuclei may only be expected in long bursts with observed gamma-ray durations \( t_s \approx 10 - 15(1 + z) \) s, while longer burst jets would be expected to consist mainly of protons. For the GRB external shock and hypernova source scenarios, the nuclei may be the heavy elements present in the stellar wind of Wolf-Rayet stars, which are thought to be the progenitors of GRBs and hypernovae. The stellar wind of Wolf-Rayet stars, especially the WO and WC subtypes, is heavily enriched with intermediate-mass nuclei, such as O, C, etc.

After they escape from their sources (see, e.g., Dermer 2007b), UHECR nuclei can also be subject to photodisintegration in intergalactic space, before arriving at Earth. An UHE nucleus with an energy \( \gtrsim 10^{19} \) eV will mainly collide with cosmic infrared background (CIB) photons. Using new constraints on CIB data, it has been found that UHE iron nuclei with energies \( \lesssim 10^{20} \) eV have a mean free path of \( \gtrsim 500 \) Mpc, and that this mean free path increases rapidly as the energy of the nucleus decreases (Hooper et al. 2007; Stecker & Salamon 1999). Thus, UHE nuclei can in principle originate from sources at cosmological distance, such as GRBs and hypernovae. The main question appears to be whether they survive the environment of their original sources. For this reason, most of the present paper has been devoted to a quantitative study of the survival of UHE nuclei in these sources. We find that:

1. In GRB internal shocks, heavy nuclei can survive in the sources if the internal shock radius and/or the Lorentz factor of the relativistic jets are relatively large, as given by equation (11). Thus, one might expect acceleration of Fe nuclei to be more favored in bursts with smoother, longer variability timescale light curves. For a smaller internal shock radius, the photodisintegration process due to prompt X-ray photons may become optically thick. However, if a synchrotron self-absorption break is present above several keV in the photon spectrum, the reduced number density of X-ray photons will lower the photodisintegration optical depth accordingly. In general, compared to heavy nuclei, the UHE intermediate-mass nuclei find it easier to survive photodisintegration in internal (and external) shocks.

2. In GRB external shocks, due to the much larger dissipation radii compared to internal shocks, UHE nuclei can easily survive in the sources. We have also calculated the evolution of the photodisintegration optical depth during the afterglow phase for both an ISM external medium and a stellar wind medium. The results show that UHE nuclei can survive in the afterglow shock for a relatively long time in both cases.

3. In the hypernova remnant acceleration scenario, UHE nuclei can survive in the sources, except in the early short period of time (\( \lesssim 100 \) days; for typical parameters, see eq. [31]). In this early short period, however, only a very small amount of ejecta energy has been converted to UHECRs, so their contribution to the UHECR flux is negligible. Most of the UHECRs originate from the blast wave ejection with \( \Gamma \beta \lesssim 0.5 \), which decelerate after \( \sim 1000 \) days, typically. UHE nuclei accelerated during this time are safe from photodisintegration.

Earlier calculations on the photodisintegration of UHE nuclei in GRB internal shocks by Anchordoqui et al. (2008) consider only one specific set of typical parameters for the internal shock radius and the relativistic Lorentz factor. We have improved this by considering the whole parameter space and working out the constraints on the physical conditions under which UHE heavy nuclei can survive in internal shocks. We have also considered the effect of a self-absorption break in the prompt emission on the photodisintegration problem. Besides internal shocks, we have explored the photodisintegration problem of UHE nuclei in GRB external shocks and hypernova remnants.

In summary, we have suggested possible scenarios for the injection of heavy nuclei into the acceleration zones of extragalactic sources such as GRBs and hypernovae. We found that, if heavy nuclei are accelerated in these sources, they will survive the threat of photodisintegration under fairly general conditions for the case of GRB external shocks, and for hypernovae. They could survive also in GRB internal shocks, if the latter occur at relatively large radii and/or the bulk Lorentz factors are large. On the other hand, for small shock radii and/or smaller bulk Lorentz factors, a pure proton UHECR composition would be favored if the self-absorption break in the photon spectrum is not high. Since the instability-induced entrainment process of heavy nuclei in internal shocks is currently not well known, we stress that the expected fraction of heavy nuclei injected into the internal shock acceleration process is uncertain. A significant fraction is plausible, but if the entrainment is inefficient, we would in any case expect a proton-dominated composition in GRB internal shock models. In GRB external shocks and hypernova shocks, on the other hand, the abundance of heavy nuclei is dependent on the external medium or the stellar outer envelope and wind enrichment fraction.
As we were completing our work, the Pierre Auger Collaboration (2007) reported a plausible correlation between UHECR at energies above $6 \times 10^{19}$ eV, assuming that they are protons, and AGNs within 75 Mpc selected from a particular catalog. Although this result is statistically significant at the 2.8 $\sigma$ level, as the authors themselves stress, it does not rule out a possible origin in sources, e.g., galaxies, which have a distribution similar to that of the AGNs. There is an unresolved tension between the above Auger spatial correlation analysis suggesting protons, and previous Auger results on maximum shower elongations $X_{\text{max}}$ suggesting a significant heavy element component for UHECR in the same energy range. The AGN hypothesis for the UHECRs has been questioned in the analysis of Gorbunov et al. (2007). In this paper, we have investigated the conditions under which UHECRs originating from GRBs or hypernovae could contain a significant fraction of heavy nuclei, as well as the conditions under which these UHECRs would be expected to be mainly protons. If any of the observed UHECRs were heavy nuclei, they could reach the Earth from distances beyond 100 Mpc, and the number of galaxies which can host a GRB or hypernova increases significantly with distance. Within the proton-inspired $3.2^{+2}_{-1}$ circle containing an AGN in the Auger correlation analysis, there would be many more galaxies that could host a GRB or a hypernova. In the case of heavy nuclei, the deflection angles are larger (e.g., for oxygen with energy $\geq 6 \times 10^{19}$ eV it is $\leq 16^\circ$; Sommers 2007), which would contain even more galaxies. While a study of the angular correlations is beyond the scope of this paper, our results are compatible with and relevant for both the Auger elongation studies and spatial correlation studies, providing constraints which can be used in future analyses.

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