THE SURVIVAL OF NUCLEI IN JETS ASSOCIATED WITH CORE-COLLAPSE SUPERNOVAE AND GAMMA-RAY BURSTS

SHUNSAKU HORIUCHI1, KOHTA MURASE1, KUNIHITO IOKA2, AND PETER MÉSZÁROS3
1 Center for Cosmology and Astro-Particle Physics, The Ohio State University, 191 West Woodruff Avenue, Columbus, OH 43210, USA
2 KEK Theory Center and the Graduate University for Advanced Studies (Sokendai), Tsukuba 305-0801, Japan
3 Department of Astronomy and Astrophysics, Department of Physics and Center for Particle Astrophysics, 525 Davey Laboratory, Pennsylvania State University, University Park, PA 16802, USA
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ABSTRACT

Heavy nuclei such as nickel-56 are synthesized in a wide range of core-collapse supernovae (CCSN), including energetic supernovae associated with gamma-ray bursts (GRBs). Recent studies suggest that jet-like outflows are a common feature of CCSN. These outflows may entrain synthesized nuclei at launch or during propagation, and provide interesting multi-messenger signals including heavy ultra-high-energy cosmic rays. Here, we investigate the destruction processes of nuclei during crossing from the stellar material into the jet material via a cocoon, and during propagation after being successfully loaded into the jet. We find that nuclei can survive for a range of jet parameters because collisional cooling is faster than spallation. While canonical high-luminosity GRB jets may contain nuclei, magnetic-dominated models or low-luminosity jets with small bulk Lorentz factors are more favorable for having a significant heavy nuclei component.

Key words: gamma-ray burst — nuclear reactions, nucleosynthesis, abundances — supernovae: general

Online-only material: color figures

1. INTRODUCTION

Recent studies have led to a canonical picture where long gamma-ray bursts (GRBs) are rare types of core-collapse supernovae (CCSN) that are accompanied by the launch of energetic relativistic jets (see, e.g., Mészáros 2006; Woosley & Bloom 2006 for reviews). While the GRB is a rare phenomenon, a significantly larger fraction of CCSN could produce jets that do not successfully produce GRBs either because of energetic or collimation reasons. Many such jets may even be choked in their progenitor envelopes (Mészáros & Waxman 2001). Observationally, spectropolarimetry shows that the degree of asymmetry in CCSN increases with time and hence to greater depths of the CCSN, suggesting an association with the central engine and possible launch of bipolar jets (Wang et al. 2001; Chornock et al. 2010). SN 1987A, the closest CCSN in modern times, shows a globally asymmetric expanding debris with an axis that roughly aligns with that of its rings (Wang et al. 2002). Cas A, one of the well-known CCSN remnants in our Galaxy, may have been accompanied by an iron-rich jet component which may help to explain the observed apparent overturn of the Fe-rich ejecta (Wheeler et al. 2008; DeLaney et al. 2010).

It is well known that nuclei can be synthesized not only in stellar nucleosynthesis but also during the CCSN, where presolar nuclei in the star are further fused into heavier nuclei via explosive nucleosynthesis. For CCSN associated with GRBs, such as SN 2003dh (GRB 030329) and SN 1998bw (GRB 980425), large kinetic energies of $\sim 10^{52}$ erg are suggested and the inferred synthesized $^{56}$Ni masses are $\sim 0.5 M_\odot$ (Iwamoto et al. 1998; Woosley et al. 1999; Woosley & Heger 2003). The large explosion energies may be due to a baryon-rich jet component driving the CCSN, which would lead to different elemental yields due to additional nucleosynthesis and/or efficient mixing and transport of stellar nuclei (Lemoine 2002; Pruet et al. 2002; Maeda et al. 2002; Maeda & Nomoto 2003). Alternatively, a large amount of $^{56}$Ni can be produced by a disk wind, where free nucleons ejected from the disk combine to form heavy nuclei (MacFadyen & Woosley 1999; Pruet et al. 2003; Surman & McLaughlin 2005). The highly relativistic jet responsible for the GRB should have a smaller baryon mass, but may still pick up and/or entrain stellar nuclei, nuclei synthesized by a wider jet, or nuclei synthesized by a disk wind.

The fate of nuclei in relativistic jets is also interesting in light of recent reports on the nuclear composition of ultra-high-energy cosmic rays (UHECRs). The origin of UHECRs remains one of the great mysteries in high-energy astrophysics (see, e.g., Hillas 2005; Blümmer et al. 2009; Beatty & Westerhoff 2009; Kotera & Olinto 2011 for reviews). At the highest energies (above $\sim 50$ EeV), the extreme energies argue for extragalactic sources (although Galactic transients may give a contribution around $\sim 10$ EeV; e.g., Calvez et al. 2010). Source candidates fall into active galactic nuclei (AGNs), including radio-loud AGNs (Biermann & Strittmatter 1987; Takahara 1990; Norman et al. 1995), radio-quiet AGNs (Pe’er et al. 2009), young AGNs (Takahara 1990; Takami & Horiuchi 2011), as well as powerful transient flares (Farrar & Gruzinov 2009; Dermer et al. 2009); GRBs (Waxman 1995; Vietri 1995; Milgrom & Usov 1995), including sub-luminous GRBs accompanying relativistic ejecta in some form (Murase et al. 2006, 2008; Wang et al. 2007; Chakraborti et al. 2011; Liu & Wang 2011); the formation of rapidly rotating strongly magnetized protomagnets (Arons 2003; Murase et al. 2009; Metzger et al. 2011; Kotera 2011; Fang et al. 2012); and galaxy cluster shocks (Norman et al. 1995; Kang et al. 1996; Inoue et al. 2007). Non-astrophysical sources include top-down scenarios where the UHECR arise from decays of massive relics from the early universe (e.g., Berezhinsky et al. 1997).

The composition of UHECRs is observationally inferred by measuring composition-dependent quantities of the showers made as UHECRs enter Earth’s atmosphere. The Pierre Auger Observatory (PAO) finds that the average shower depth at shower maximum, $X_{\text{max}}$, and its rms variation $\sigma(X_{\text{max}})$, suggest the composition becoming increasingly dominated by heavy nuclei above the “ankle” of $10^{18.5}$ eV (Abraham et al. 2010a; Abreu et al. 2011a). Note that the PAO results have not been verified by the High Resolution Fly’s Eye experiment (Abbasi et al.
2010b) and the Telescope Array experiment (e.g., Tsunesada et al. 2011), although the latter experiments observe a different hemisphere and their statistics are lower. In addition, these experiments do not confirm the directional anisotropies reported by PAO (Abraham et al. 2007; Abreu et al. 2010b; Abbasi et al. 2008, 2010a). Also, the PAO indicators have been shown to be inconsistent with each other (Shaham & Piran 2012) and that they may not reflect a composition change given the current uncertainty in shower interaction physics (Wilk & Włodarczyk 2011). More observational and theoretical studies are required to settle these and other remaining issues.

If UHECRs are indeed composed of heavy nuclei, this provides important implications for UHECR sources. In the AGN and galaxy cluster shock origin scenarios, the matter being accelerated originates from the intergalactic medium, so that the mass fraction of Fe and heavier nuclei is small: ~10^{-3} for solar metallicity. Although the dominance of nuclei in UHECRs may be explained by a rigidity-dependent acceleration mechanism, where the maximum Fe energy would be $E_{F_{\text{max}}} = 26$ times higher than those of protons, this requires the maximum accelerated energies in all contributing sources to be somewhat fine-tuned such that $E_{F_{\text{max}}} \approx 10^{19}$ eV. Also, such a scenario predicts a larger proton-to-nuclei ratio at low energies than is actually observed (Lemoine & Waxman 2009; Abreu et al. 2011b). An alternate scenario is that the environment supplying the injected particles is enriched in nuclei. The abundance of nuclei inferred from $X_{\text{max}}$ and $\text{rms}(X_{\text{max}})$ require high nuclei abundances (see, e.g., Anchordoqui et al. 2007; Taylor et al. 2011 and references therein). Massive-star origins, including GRBs, CCSN with relativistic ejecta, and magnetars would be attractive sources in this sense because as described above the environment contains large fractions of intermediate or heavy nuclei. Additionally, it has been shown that once loaded, nuclei may be accelerated to ultra-high energies and successfully survive in the dissipation regions of jets, including both classical high-luminosity GRBs (HL GRBs) and sub-luminous GRBs (Murase et al. 2008; Wang et al. 2008).

In this paper, we investigate the origins and survival of nuclei as the jet is still inside the star. More specifically, we first discuss locations where nuclei may enter the jet medium and investigate whether nuclei can survive in each of these locations. Second, we investigate whether nuclei that have successfully made their way into the jet survive during the jet propagation through the progenitor star. These issues are different from previous works that mainly focused on the survival of UHECR nuclei in the emission regions (where the GRB occurs, typically outside the star). Note that UHECR acceleration is not expected inside the star even though our study would be useful for UHECR sources, and we mainly consider the survival of low-energy nuclei.

The paper is organized as follows. In Section 2, we discuss sources of nuclei and discuss conditions for nuclei survival in each of them. In Section 3, we consider processes in the jet that may lead to nuclei disintegration and discuss conditions under which nuclei survive. Finally, we finish with discussions and conclusions in Section 4. We express quantities as $Q_{c} = Q/10^{5}$ in cgs units.

2. NUCLEI ORIGINS AND SURVIVAL

We first discuss three potential sources of nuclei in the context of GRBs. First, loading at the jet base, where a small amount of baryons may be loaded into the jet (the jet is assumed to be radiation or magnetically dominated). Depending on the jet parameters, nuclei survive or are disintegrated into free nucleons. Second, nucleosynthesis in the outflow itself, which occurs if nuclei are disintegrated into free nucleons at initial loading. Finally, entrainment, which is the loading of surrounding nuclei into the jet during jet propagation.

2.1. Nucleus Loading at the Jet Base

The progenitors of GRBs are massive stars with Fe cores of mass approximately 1–2 $M_{\odot}$ that extend to a few $\times 10^{6}$ cm. It is thought that a few seconds after the onset of collapse, a compact object (either a neutron star or a black hole) forms at the center, surrounded by an equatorial disk supported by centrifugal forces. By some currently unconfirmed mechanism, an energy reservoir is tapped and energy is released near the collapsing core. In the “collapsar” model, the energy derives from the gravitational energy of rapid accretion, and neutrinos play the role of energy transport (Woosley 1993; MacFadyen & Woosley 1999). In this case, the energy is deposited dominantly as radiation energy. Alternatively, the jet may be dominated by a magnetic component, as, for example, in scenarios where the jet magnetic field threads the black hole event horizon and the power derives from the black hole spin energy (Blandford & Znajek 1977), where the jet magnetic field threads the surface of the accretion disk (Blandford & Payne 1982; Proga et al. 2003), or where the power derives from the spin-down energy loss of a central magnetic proto-neutron star (Usos 1992).

We consider a jet with baryonic matter injected at the rate $\dot{M}_{0}$, radiation energy at the rate $\dot{E}_{\text{rad},0} \gg \dot{M}_{0} c^{2}$, and magnetic energy at the rate $\dot{E}_{\text{mag},0}$ (we will use collimation-corrected values, not the isotropic equivalent) at a radius of $r_{0}$. The jet subsequently adiabatically expands as its radiation and/or magnetic energies are gradually converted to bulk kinetic energy. The maximum bulk Lorentz factor is set by $\eta = L_{\text{ke}}/(M c^{2})$, where $L_{\text{ke}} = L_{\text{rad},0} + L_{\text{mag},0}$ is the kinetic luminosity after the end of the jet bulk acceleration phase. We defined a “fireball” as a jet where radiation dominates, i.e., $L_{\text{rad},0} \gg L_{\text{mag},0}$, and a magnetic outflow as the opposite with $L_{\text{mag},0} \gg L_{\text{rad},0}$.

The large values of $\eta \gtrsim 100$ required in GRBs (Ruderman 1975) necessitate a jet with small baryon loading. The quantitative predictions about where and how baryon loading occurs are tentative at best, but any baryons must come from the surrounding material. The composition of the pre-supernova stellar core is dominantly nuclei (bottom panel of Figure 1). At small radii thus can be significantly altered by the collapse. For example, if a supernova shock is launched prior to the jet, nuclei out to a distance of $\sim 10^{3} - 10^{4}$ cm would be disintegrated by the supernova shock. At later times, explosive nucleosynthesis can alter the composition. For example, Fryer et al. (2006) use one-dimensional simulations to demonstrate with that up to $\sim 1 M_{\odot}$ of $^{56}\text{Ni}$ may be synthesized, depending on whether the core collapse proceeds to a black hole directly or via fallback, and also on the explosion energy. The detailed composition of the surroundings therefore depends on the GRB scenario and the timing of jet launch relative to core collapse. In either case, it is expected that there are abundant nuclei especially for radii larger than $\sim 10^{7}$ cm.

Any nuclei that are loaded into the jet are destroyed if it collides with particles or photons with energies exceeding the nuclear binding energy ~10 MeV in the nuclei rest frame. We consider a jet of luminosity $L_{\text{rad},0}$ injected at a radius $r_{0}$ with $\approx 2.10^{19}$ eV.

4 We caution that in magnetic-dominated jets such definitions may not strictly hold because of the model-dependent conversion of magnetic energy to kinetic energy; see Section 2.3.1.
Also, ... radiation fireball, $L_{rad,0}$ corresponds to $L_{ke}$ and the GRB luminosity by Equation (34). In the case of a magnetic-dominated outflow, $L_{rad,0}$ may be significantly smaller for a given observed GRB luminosity and nuclei loading would be easier. The bottom panel shows the dominant pre-supernova stellar abundances taken from the rotating 20 $M_\odot$ pre-supernova progenitor model E20 of Heger et al. (2000). The abundances are illustrative: depending on the timing of jet launch, nuclei at small radii will be disintegrated by a supernova shock or synthesized by explosive nucleosynthesis before the jet propagates through (see the text).

(A color version of this figure is available in the online journal.)

initial Lorentz factor $\Gamma_0 = 1$. Its temperature is related to the radiation energy density,

$$aT_0^4 = \frac{L_{rad,0}}{\Sigma_0h\Gamma_0^2c^3}, \quad (1)$$

where $a$ is the radiation constant and $\Sigma_0 = \Omega_0 r_0^2$ is the jet cross section. Typically, the jet is collimated already at injection due to rotation or magnetic geometries, and we adopt a solid angle $\Omega_0 = 0.1$ sr independent of $r_0$. This results in

$$T_0 \approx 1.3 \Omega_0^{-1/4} L_{rad,0}^{1/4} r_0^{-1/2} \text{MeV}. \quad (2)$$

Therefore, nuclei will be destroyed by the high-energy tail of the photon spectrum. The optical depth for photodisintegration is $\tau_{A\gamma} \approx n_0 \sigma_{A\gamma} r_0 / \Gamma_0$, where $n_0$ is the photon density and $\sigma_{A\gamma}$ is the photodisintegration cross section which peaks at approximately $10^{-25}$ cm$^2$ (e.g., Murase et al. 2008). We require that $\tau_{A\gamma} < 1$ for the survival of nuclei during loading. Similarly for spallation we require that $\tau_{sp} \approx n_0 \sigma_{sp} r_0 / \Gamma_0 \ll 1$, where $n_0$ is the comoving ion density,

$$n_0 = \frac{L_{ke}}{\Sigma_0 h \Gamma_0^2 A m_p c^3}, \quad (3)$$

and $\bar{A}$ is the average mass number of the baryons in the jet. Also, $\sigma_{sp} = \sigma_0 A^{2/3}$ with $\sigma_0 \approx 3 \times 10^{-26}$ cm$^2$ and $A = 56$ is the mass number of the injected nuclei of interest (assumed Fe).

In Figure 1, the photodisintegration and spallation optical depths are shown for three values of $L_{rad,0}$. We adopt a thermal photon spectrum of temperature $T_0$, and assume that the jet ions have a Maxwell–Boltzmann velocity distribution with temperature equivalent to the photon temperature. Photodisintegration is more important than spallation because the larger photon density compensates the slightly smaller peak photodisintegration cross section. Note that in Figure 1 we adopt $A = 1$, which gives the largest spallation target density possible.

As expected, for a GRB fireball of typical luminosities ($L_{rad,0} = 10^{49} - 10^{51}$ erg s$^{-1}$), the high $T_0$ results in nuclei being photodisintegrated at loading. However, since photodisintegration occurs with the exponential tail of the photon spectrum, nuclei survival is highly sensitive on $T_0$ and hence GRB parameters. For example, nuclei survival is possible for $r_0$ greater than $10^8 - 10^9$ cm. Values beyond $10^9$ cm are somewhat large but plausible if the jet is powered by an accretion disk. However, in this case it may be difficult to achieve short variable timescales on the order of ms observed in some GRBs, unless they are imprinted by instabilities during jet propagation. Figure 1 also demonstrates that nuclei survival is more easily possible for sub-luminous GRBs. We quantify these constraints in Figure 2.

It is easier to initially load nuclei in magnetically dominated jets ($L_{mag,0} \gg L_{rad,0}$) because $L_{rad,0}$ is smaller for a given observed GRB luminosity and $T_0$ is correspondingly smaller. Both photodisintegration and spallation become less effective.

2.2. Nucleosynthesis by Outflows

The fireball is composed of free nucleons if the photodisintegration optical depth at jet launch is larger than unity. As the jet expands and cools, nucleons combine to form nuclei. The freezeout composition of HL GRB jets has been investigated by various authors, and strongly depends on the radiation energy content of the jet. For a canonical GRB fireball of initial temperature $\sim 1$ MeV, the entropy per baryon is necessarily high, $S \gtrsim 10^5 k_b$ nucleon$^{-1}$. Free nucleons combine to form $\alpha$ particles only once the deuterium bottleneck has been broken, which
occurs at lower jet densities for higher entropies where the further processing into carbon and higher nuclei are not fast enough compared to the expansion timescale. As a result, few elements heavier than He are formed, similar to big bang nucleosynthesis (Lemoine 2002; Pruet et al. 2002; Beloborodov 2003).

The situation is dramatically different for magnetically dominated jets, where most of the energy is stored in magnetic energy and the jet entropy is necessarily lower, \( S \sim 10–300 \, k_b \) nucleon\(^{-1}\). Under such conditions, \( \alpha \) recombination occurs at higher densities such that heavier elements can be formed efficiently via the triple-\( \alpha \) process and subsequent \( \alpha \) captures. Focusing on the magnetic jets of a millisecond protomagnetar central engine, Metzger et al. (2011) showed that the jet composition may indeed be dominated by heavy nuclei. Also, depending on how neutron-rich the jet matter is (Metzger et al. 2008), neutron-catalyzed \( \alpha \) formation and neutron captures open the path for the rapid synthesis of Fe-peak and heavier \( n \)-capture nuclei (e.g., \( A \gtrsim 90 \)). Similar nucleosynthesis is also realized in more baryon-rich jets with lower entropies (Inoue et al. 2003). Although the baryonic jets cannot produce the classical GRB phenomenom, it can be relevant for sub-luminous GRBs and hypernovae where \( \eta \) can be much lower than \( \eta \sim 100–1000 \).

When the central engine consists of a black hole and accretion disk system, hot outflows from the disk are expected, where the disk wind with \( S \sim 10–100 \, k_b \) nucleon\(^{-1}\) can account for the amount of \( ^{56}\text{Ni} \) observed in hypernovae associated with GRBs (e.g., Surman & McLaughlin 2005). We discuss the entrainment of external nuclei next.

\subsection*{2.3. Nucleus Entrainment During Propagation}

As the jet propagates through the star, it can pick up baryons from the stellar core, surrounding wider jet, or the disk wind environment. Since these environments can be nuclei-rich, we discuss how nuclei may survive during the entrainment process. We will assume a proton-dominated jet (\( A = 1 \)) for this purpose. For our canonical jet, we consider a classical HL GRB jet, fixing \( \Gamma_0 = 1 \) and \( \Omega_0 = 0.1 \) sr. Further, we adopt \( L_{\text{ke}} = 10^{50} \) erg s\(^{-1}\), \( \eta = 100 \), and \( r_0 = 10^3 \) cm, and show dependencies where applicable. For illustration purposes, we adopt the rotating 20 \( M_\odot \) pre-supernova progenitor model E20 of Heger et al. (2000), whose density profile can be well approximated by \( \rho_\nu(r) = 1.7 \times 10^5 (r/10^9 \text{ cm})^{-3} \, g \text{ cm}^{-3} \) with \( n \sim 3 \).

\subsubsection*{2.3.1. Cocoon Properties}

After the onset of collapse, the lack of centrifugal force along the rotational axis leads to free mass infall and the formation of a funnel region where conditions are favorable for the launch of a jet (Woosley 1993; MacFadyen & Woosley 1999). In magnetic-dominated jets, the outflow initially flows along open field lines. However, self-collimation via magnetic hoop stress fails in relativistic flows (e.g., Sakurai 1985; Bucciantini et al. 2006), and recent works show how interactions with the star may redirect the flow toward the poles and produce relativistic bipolar jets along the rotational axis (Bucciantini et al. 2007, 2008; Komissarov & Barkov 2007; Uzdensky & MacFadyen 2007). We model the jet cross section as

\[ \Sigma(r) = \Sigma_0 \left( \frac{r}{r_0} \right)^{\xi}, \]

where, for example, \( \xi = 2 \) corresponds to a conical (or spherical symmetric) jet, and \( \xi = 1 \) corresponds to a funnel-like jet (e.g., Mészáros & Rees 2001).

After injection, the GRB jet accelerates as its radiation or magnetic energy is converted to bulk kinetic energy. For a radiative fireball expanding adiabatically with cross section \( \Sigma \), energy conservation yields (see also, e.g., Toma et al. 2007; Ioka et al. 2011)

\[ \Gamma_j(r) = \begin{cases} \Gamma_0 (r/r_0)^{5/2} / \eta & r < r_{\text{sat}} \\ 1 & r > r_{\text{sat}} \end{cases}, \]

where \( r_{\text{sat}} = \eta^{2/5} r_0 \) is the saturation radius where the terminal Lorentz factor is reached. Often \( r < r_{\text{sat}} \) is called the jet or bulk acceleration phase and \( r > r_{\text{sat}} \) is the jet coasting phase.

The relativistic jet slows down abruptly in a narrow layer at the head of the jet where it comes into contact with the overlying stellar material. Due to the large ram pressure experienced, a strong reverse shock forms and the jet head is decelerated to a sub-relativistic velocity \( \beta_h < 1 \). Further upstream, a forward shock forms that shock heats the stellar material. The jet head velocity is set by balancing the jet pressures applied on the forward and reverse shocks. We use the analytic approximations of Matzner (2003),

\[ \beta_h = \beta_j \frac{1}{1 + \tilde{L}^{-1/2}}. \]

In the limit \( \tilde{L} \ll 1 \), which corresponds to a relativistic reverse shock (as viewed in the jet frame) as is generally the case for jet propagation in the star, Equation (6) can be linearized, and

\[ \beta_h \approx \tilde{L}^{1/2} \propto L_{\text{ke}}^{1/2} n^{(1-\xi)/2} r_0^{(\xi-2)/2}. \]

We call the shocked stellar and shocked jet material collectively as the jet head. At first, the jet head pressure is smaller than the surrounding stellar pressure and the jet head remains pressure confined. The jet head pressure is \( P_h \approx U_h / 3 \), where \( U_h \approx 4 \Gamma_j^2 n_p m_p c^2 \) is the jet head energy density and \( n_p \) is the proton density, i.e., the ion density

\[ n_i(r) = \frac{L_{\text{ke}}}{\Sigma(r) n \Gamma_j A m_p c^3}, \]

with \( \tilde{A} = 1 \). Hence, \( P_h \) falls as \( \propto r^{-\xi/2} \) during the bulk acceleration phase and \( \propto r^{-\tilde{A}} \) during the coasting phase, while the stellar pressure falls as \( P_\nu \propto \rho_n^{4/3} \propto r^{-4\xi/3} \). Therefore, the jet head pressure eventually overtakes the stellar pressure and starts to overflow. We call the radius at which the pressures become equal \( r_c \); cocoon formation occurs for \( r > r_c \). For typical parameters, \( r_c \) is in the range \( 10^8–10^9 \) cm and cocoon formation occurs during the bulk acceleration phase (Mészáros & Rees 2001).

For magnetic-dominated outflows, the conversion of magnetic energy to kinetic energy is prolonged since only part of the magnetic luminosity gets converted directly into kinetic luminosity, the other part being converted to kinetic energy in a two-stage process through thermal energy (e.g., Mészáros & Rees 2011). Furthermore, the conversion is not efficient in unconfined, time-stationary outflows in ideal MHD (e.g., Goldreich & Julian 1970; Bogovalov & Tsinganos 1999). Therefore, models...
for full jet acceleration employing a combination of differential collimation, time variability, or violations of ideal MHD have been studied. These result in $\Gamma$, that increases roughly as $\propto r^{1/3}$ (e.g., Drenkhahn 2002; Granot et al. 2011; Metzger et al. 2011) or $\propto r^{1/2}$ (e.g., McKinney & Uzdensky 2012). For illustration, we adopt the calculations of McKinney & Uzdensky (2012) who find the jet dynamics well fit by $\Gamma_j \propto r^{1/2}$ and $\Sigma \propto r^{5/4}$ during jet propagation in the star (see their Figure 5). In this case the jet head pressure falls as $\propto r^{-3/4}$, so cocoon formation is again inevitable, occurring in the range $10^8–10^9$ cm. Note that our adopted magnetic-dominated jet has properties—such as particle density, energy density, bulk acceleration, and so on—in between our conical and funnel jets.

### 2.3.2. Survival of Nuclei in the Cocoon

The freshly formed cocoon consists of an inner region composed of shocked jet material and an outer region composed of shocked stellar material. The two are separated by a contact discontinuity. The contact discontinuity is dynamically unstable and studies find that it remains on the order of seconds (e.g., Mizuta & Aloy 2009). After that, the outer and inner cocoon material mix. We discuss how stellar nuclei may survive until this last phase of cocoon evolution. Note that the discussions below will also apply to magnetic outflows if they interact with the progenitor to produce a cocoon.

Nuclei in the star ahead of the jet are first shocked by the forward shock approaching with velocity $-\beta_h$. Spallation occurs when the energy of a collision between nuclei and jet head protons exceeds the nuclear binding energy $\approx 10$ MeV in the nuclei rest frame. Therefore, external nuclei survive if the jet head velocity is less than $\beta_{sp}$ in the stellar frame. The jet head bulk velocity for a conical ($\xi = 2$) jet,

$$\beta_h^{(2)} \sim 0.01 L_{ke,50}^{1/2} r_c^{1/2},$$

typically exceeds $\beta_{sp}$ for radii close to $10^{11}$ cm. For more collimated jets, $\beta_{sp}$ is reached at smaller radii. For example, the jet head velocity of a $\xi = 1$ jet of the same jet parameters is $\beta_h^{(1)} \sim 0.1 L_{ke,50}^{1/2} r_c^{1/2} r_0^{-1/2}$. Also note that $\beta_h$ grows faster with radius and the saturation radius is larger. Therefore, this channel for nucleon loading is limited to jet head radii less than $10^8–10^9$ cm depending on jet parameters.

Nuclei can also enter the cocoon directly through its boundary with stellar matter. The cocoon is overpressured and expands into the star at velocity $\beta_c$ given by pressure equilibrium, $\rho_s c^2 \beta_c^2 = P_e$, where $P_e = E_c / (3V_c)$ is the radiation dominated cocoon pressure and $E_c \approx L_{ke}(t-r/c)$ is the total energy deposited in the cocoon (Begelman & Cioffi 1989; Matzner 2003). We approximate the cocoon as a cylinder of height $r$ and base length $x_c \approx c \beta_c t$, so that its volume is $V_c = \pi c^2 \beta_c^2 t^2 r$. The cocoon expansion velocity is then

$$\beta_c \approx \left( \frac{P_e}{\rho_c c^2} \right)^{1/2} \propto L_{ke}^{3/8} r^{(3n-\xi-4)/8} r_0^{(\xi-2)/8},$$

which is significantly slower than the head velocity and depends weakly on jet parameters. For example, even for a $\xi = 1$ jet with our nominal jet parameters (i.e., $L_{ke} = 10^{50}$ erg s$^{-1}$, $n = 100$, and $r_0 = 10^7$ cm), the cocoon expands at only $\beta_{c}^{(1)} \sim 0.04 r_0^{1/2}$. Furthermore, the flow rate of material into the cocoon through the cocoon boundary can be substantial. The flow rate scales as the product of the expansion velocity and surface area. For the jet head this is $\sim \pi x_c^2 c \beta_h$. For the cocoon boundary this is $\sim 2\pi x_c r_c \beta_h$, which can be rewritten as $2\pi x_c^2 c \beta_h$ using the approximations $r \approx \beta_h t$ and $x_c \approx c \beta_c t$. Thus, the flow rates through the jet head and cocoon boundary are likely at least comparable.

Once in the outer cocoon, nuclei must survive photodisintegration. The photon temperature in the bulk acceleration phase is

$$T_h \approx \left( \frac{e_c U_h}{a} \right)^{1/4} \propto e_c^{1/4} L_{ke}^{1/4} r^{-8/7} r_0^{-1/4} t_0^{(\xi-4)/8},$$

and for a conical ($\xi = 2$) jet this yields

$$T_h^{(2)} \sim 100 e_c^{-1/4} L_{ke,50}^{1/4} r_c^{-1/4} r_0^{-1/4} t_0^{-1/4} \text{keV},$$

where for generality we adopt a parameter $e_c \lesssim 1$ to characterize the fraction of the jet head internal energy that is radiated by electrons. Throughout, we adopt $e_c = 0.1$ and show dependencies on $e_c$. We quote results for the bulk acceleration phase because the GRB jet is typically in the acceleration phase when it advances through the stellar core, unless it is a baryon-rich jet with a small saturation radius. Note that for a $\xi = 1$ jet of the same parameters, $T_h^{(1)} \sim 160$ keV. Thus, temperatures are insufficient for rapid photodisintegration. Also, the cocoon would be turbulent, with flow speeds that may reach the sound velocity $c/\sqrt{3}$. Nuclei being carried by such flows will exceed the spallation energy threshold. As we detail in Section 2.3.3, only relativistic nuclei above a critical energy are spalled because lower energy nuclei are expected to thermalize rapidly.

Next, we must check that nuclei survive potential collisions with protons in the inner cocoon. In particular, jet protons may acquire relativistic random velocities at the reverse shock. The protons could be trapped in the fireball by strong magnetic fields that are either generated in situ or advected from the central engine, and at the reverse shock their directions could be isotropized so that they maintain their velocities comparable to $\Gamma_j$ in magnitude (e.g., see discussions in Ioka 2010). In this case, the relativistic protons can cause spallation reactions. However, the relativistic protons typically lose energy very rapidly. For example, relativistic protons lose energy by $\pi$-production at an energy loss rate

$$v_{pp} \sim (4\Gamma_j n_p) a_0 c \propto L_{ke} r^{-\xi} r_0^{\xi-2},$$

where the quantity in brackets is the jet head particle density and $a_0$ is again $3 \times 10^{-20}$ cm$^{-2}$. For our canonical jet, $v_{pp}^{(2)} \sim 2 \times 10^8 r_0^{-2} s^{-1}$ and $v_{pp}^{(1)} \sim 1 \times 10^{10} r_0^{-1} s^{-1}$ for $\xi = 2$ and $\xi = 1$, respectively. $\pi$-production reduces the proton kinetic energy down to around 70 MeV. Protons also lose energy by processes such as $e^\pm$-production and Coulomb interactions which are particularly important to reduce the proton kinetic energy further. Since electrons rapidly lose their energy by Compton scattering and Bremsstrahlung emission, the electrons thermalize and their temperature is lower than protons. Thus, $\sim 70$ MeV protons lose energy to electrons at a rate (see also Section 2.3.3)

$$v_{pe} \sim 32 \sqrt{\eta} n_p q_4^4 \ln \Lambda_{pe} \propto L_{ke} r^{-\xi} r_0^{\xi-2},$$

where $\ln \Lambda \sim 10$ is the Coulomb logarithm, $v_e$ is the electron velocity,$^5$ and $n_p = n_e$ if the only electrons in the jet are those

$^5$ We assume that $T_e$ is equivalent to the radiation temperature $T_r$ because of Compton scattering. When the temperature is $T_r \lesssim 10^8$ eV, $v_e \approx (3T_r/m_e c)^{1/2}$ and the $v_{pe}$ dependence changes accordingly.
associated with the ions; if $e^\pm$ pair production in the jet head can increase the proton cooling rate (see Section 3.2). For our canonical GRB jet, this yields $v_{pe}^{2(1)} \sim 1 \times 10^9 \, r_0^{-2} \, \text{s}^{-1}$ and $v_{pe}^{2(1)} \sim 7 \times 10^9 \, r_0^{-1} \, \text{s}^{-1}$ for $\xi = 2$ and $\xi = 1$, respectively.

The short cooling timescales imply that relativistic protons occupy only a very thin region around the reverse shock. The proton cooling timescale is also much shorter than the lifetime of the contact discontinuity separating the outer and inner cocoons. Thus, relativistic protons would have lost much of their energy by the time nuclei in the outer cocoon come into contact with jet protons in the inner cocoon. Note that this conclusion does not break down for jets with extremely large $\eta$, where the optical depth to $pp$ collisions can be below unity and protons will not thermalize (Ioka 2010).

Finally, nuclei must survive in the cocoon. The cocoon temperature is determined by the energy density of the cocoon. We estimate the cocoon temperature to be

\[ T_c = \left( \frac{\epsilon_e E_e}{\nu_{e'c}} \right)^{1/4} \sim \epsilon_e^{1/4} L_{ky}^{1/4} r^{-1/4} n_{e'c}^{1/2} t_{0}^{1/4} \, , \quad (16) \]

where $\epsilon_e = 0.1$ as before gives $T_c^{2(1)} \sim 80 \, r_0^{-9/16} \, \text{eV}$ and $T_c^{2(1)} \sim 100 \, r_0^{-1/2} \, \text{keV}$ for $\xi = 2$ and $\xi = 1$, respectively. These are too low for significant photodisintegration.

In conclusion, external nuclei enter the cocoon through the jet head and the boundary between the cocoon and the star. Those entering through the jet head survive provided the jet head velocity is slower than $\beta_{np}$. This is usually satisfied for $r \lesssim 10^{11} \, \text{cm}$, but for more penetrating jets can be as limited as $r \lesssim 10^9 \, \text{cm}$. Those entering through the cocoon boundary survive spallation. Once in the cocoon, nuclei survive both spallation and photodisintegration.

2.3.3. Survival of Nuclei during Entrainment

A classical HL GRB jet is likely to be highly relativistic at its central cross section, moving with bulk Lorentz factor $\Gamma_j$. Surrounding this is a transition layer where the velocity decreases from relativistic to non-relativistic values. Outside the transition layer lies the non-relativistic nuclei-rich cocoon. The growth time of shear-driven instabilities at the transition layer is much smaller than the duration of the jet, and the transition layer likely contains rapid fluctuations in thermodynamic quantities (Alford et al. 2002). Here, we discuss whether nuclei are destroyed when nuclei cross into the jet through such transition layers.

When a nucleus in the cocoon moves into the jet plasma, it has an extremely short thermal relaxation time corresponding to $\sim 1/(v_{Ap} + v_{Ae})$, where $v_{Ap}$ and $v_{Ae}$ are the energy loss rates on jet protons and electrons, respectively. These can be written as $v = 2v_{sg} - v_{l} - v_{t}$, where $v_{sg}$, $v_{l}$, and $v_{t}$ are the slowing down rate (or momentum loss rate), the pitch-angle diffusion rate, and the parallel velocity diffusion rate, respectively. The pitch angle and parallel velocity diffusion refer to the perpendicular and parallel velocity components of test particles spreading in velocity space through multiple Coulomb scatterings. For sufficiently energetic test particles, $v_{l}$ and $v_{t}$ are smaller than $v_{sg}$; in other words, the overall slowing down of a test particle is more significant than its diffusion in velocity space. We may thus approximate $v \approx 2v_{sg}$. The slowing down rate is defined $v_{sg} = -(\Delta \nu) / \nu$, where $\nu$ is the particle velocity and $v_{t}$ is in the direction of the nucleus motion (e.g., Spitzer 1956), and

\[ v_{sg,Ae} = \frac{(1 + \nu A m_A m_e) A_D G(v_{Ae}/v_{e})}{\nu A v_{e}^2} \, , \quad (17) \]

where $v_A$ is the nuclei velocity and $\gamma_A$ is its Lorentz factor, $A_D$ and $G(y)$ are

\[ A_D = \frac{8 \pi n_e q^4 Z_A^2 Z_e^2 \ln \Lambda_{eA}}{\gamma_A^2 m_A^2} \, , \quad (18) \]

\[ G(y) = \frac{\Phi(y) - y \Phi'(y)}{2 y^2} \, , \quad (19) \]

and $\Phi(y)$ is the usual error function. This is the energy loss rate on jet electrons; the same expression with $e$ replaced by $p$ applies for energy loss on jet protons. The expression becomes inaccurate when $v_{Ae}/v_e \gtrsim \ln \Lambda_{eA}$ because terms ignored in its derivation become important (Spitzer 1956). However, in our case the jet electron velocity is already close to $c$ and this is not a serious concern. Indeed, for large $y$ we can approximate $G(y) \to 1/(2y^2)$ and we find that $\nu$ agrees numerically with the well-documented energy loss rate of relativistic cosmic rays propagating through fully ionized plasmas (see, e.g., Equation (5.3.40) in Section 5.3.8.1 of Schlickeiser 2002). Note that for $v_{Ae} < v_e$, $G(y) \approx (2y)/(3\sqrt{\pi})$ and one obtains Equation (15). For reference, the expressions for the pitch-angle diffusion and the parallel velocity diffusion rates are (Spitzer 1956)

\[ v_{l,Ae} = \frac{A_D (\Phi(v_{Ae}/v_e) - G(v_{Ae}/v_{e}))}{v_A^3} \, , \quad (20) \]

\[ v_{t,Ae} = \frac{4 A_D G(v_{Ae}/v_e)}{v_A^3} \, . \quad (21) \]

These affect the energy loss rates at small nuclei energies.

In Figure 3, we show the energy loss rate $\nu = 2v_{sg} - v_{l} - v_{t}$ for a test Fe nuclei in the jet plasma at $r = 10^9 \, \text{cm}$. The jet ion component is assumed to be proton dominated ($Z_i = 1$) and the electron and proton temperatures are taken to be equivalent to the jet radiation temperature; for the bulk acceleration phase this is $T_j \simeq T_0(r/r_0)^{-\xi/2}$. For high-energy nuclei, energy loss on electrons is more important than energy loss on protons because of the faster electron velocities. We compare the energy loss rates to the spallation rate $\nu_{sp} \simeq n_p \sigma_{sp} v_{ke}$. We do not show the spallation energy loss rate because we wish to remain conservative and assume that a single spallation event affects the composition. We see that at relativistic nuclei energies, spallation dominates over Coulomb cooling rates. However, below $E_{Fe,\text{crit}} \sim 20 \, \text{GeV}$, Fe nuclei lose energy before they are spalled, even though spallation is energetically allowed. These nuclei lose energy to electrons on an exponential timescale $E_{Fe} \propto e^{-e_{\text{crit}}/m}$, where $e_{Fe,\text{crit}} = v_{Fe} (E_{Fe,\text{crit}})$. Importantly, the critical energy $E_{Fe,\text{crit}}$ depends very weakly on the radius and on GRB parameters. This is because in the limit that $v_{Ae} > v_e$, the Coulomb cooling rate can be approximated as

\[ v_{Ae} \simeq \frac{8 \pi n_e q^4 Z_A^2 Z_e^2 \ln \Lambda_{eA}}{\gamma_A m_A m_e v_A^3} \propto n_p \, , \quad (22) \]

for $Z_i = 1$. Since the spallation rate is also $\nu_{sp} \propto n_p$, the GRB dependencies largely cancel, and $E_{Fe,\text{crit}}$ is always $O(10) \, \text{GeV}$ or higher.\footnote{If $e^\pm$ pair production is important, cooling becomes more rapid and the critical energy becomes higher; see Section 3.2.} Equating $\nu_{sp}$ to $v_{Ae}$ and using the general relation $n_e = Z_i n_i$ yields the approximation

\[ E_{Fe,\text{crit}} \simeq \left[ \frac{2 \pi q^4 Z_A^2 Z_e^2 Z_i \ln \Lambda_{eA}}{\sigma_{sp} m_e / m_A} \right]^{1/2} \approx 1.4 \times 10^{10} \, \text{eV} \, , \quad (23) \]
which is close to the value in Figure 3. Note that this is a general critical energy derived from the competition between spallation and thermalization in a plasma consisting of electrons and ions of charge $Z_i$. It can also be applied to, e.g., the cocoon (Section 23.2), where $Z_i > 1$.

Let us denote the bulk velocity gradient at the cocoon–jet boundary by $d\beta/dx > 0$, where $x$ is in the transverse direction with the origin at the boundary. As a nuclei moves into the jet in the $x$-direction, collisions with jet protons become increasingly more energetic, until spallation becomes energetically possible. However, the nuclei will also tend to thermalize with the jet plasma. Once thermalized, the collision between nuclei and jet ions will lack the energy to cause spallation. Thus, if spallation is slow enough compared to the thermalization rate, nuclei will survive. Since the spallation rate depends on the collision energy, an upper limit on $d\beta/dx$ can be derived.

Consider a test nuclei moving in the $x$-direction with an initial velocity $\nu_{A,c} \sim 6 \times 10^7 T_{c,5}^{-1/2}$ cm s$^{-1}$ corresponding to $T_{c}$, and let us work in the frame of the jet immediately surrounding the test nuclei. After moving a distance $\Delta x$, if the nuclei have not yet thermalized with the surrounding jet plasma, their velocity perpendicular to $x$ is Lorentz boosted by a frame change of $c(d\beta/dx)\Delta x$. As we showed in Figure 3, nuclei survival requires the nuclei kinetic energy to be less than $E_{Fe,crit}$. This yields

$$\frac{d\beta}{dx} \Delta x < \left[1 - \left(\frac{m_A}{m_A + E_{Fe,crit}}\right)^2\right]^{1/2} \sim 0.7,$$  \quad (24)$$

where we have neglected the initial thermal kinetic energy of the nuclei as it is much smaller than $E_{Fe,crit}$. Since $E_{Fe,crit}$ depends weakly on the GRB parameters, Equation (24) does not show strong parameter dependencies either. Now, the distance

$$\Delta x \text{ governed by thermalization. Since } \nu_{A,c} \text{ falls with energy, the thermalization distance grows with energy. Larger thermalization distances result in larger frame change boosts and place stronger constraints on } d\beta/dx. \text{ So we consider the thermalization distance at } E_{Fe,crit},$$

$$\Delta x_{\text{therm}} \equiv \frac{\nu_{A,c}}{\nu_{A,c}(E_{Fe,crit})} \sim 6 T_{c,5}^{-1/2} \nu_{A,c,7}^{-1} \text{ cm}, \quad (25)$$

where we have assumed that the energetic nuclei move in a straight line (in the $x$ direction) with velocity $\nu_{A,c}$. This is an upper limit, unless nuclei are carried by e.g., bulk flows faster than $\nu_{A,c}$. From Equations (24) and (25) we can derive an upper limit on the $d\beta/dx$ required for nuclei survival. In other words, if the velocity gradient is steeper, the collision between nuclei and jet ions will become sufficiently energetic in the course of the nuclei moving a thermalization distance and spallation will occur.

In Figure 4, we show the upper limits on $d\beta/dx$ for various GRB jet parameters all for a $\xi = 2$ jet. The limits become more stringent with increasing radius, jet Lorentz factor, and inversely as the jet kinetic luminosity. These make sense, since these provide smaller target densities for nuclei to thermalize with and thus must yield stricter upper limits. For the same reason, the upper limits for a $\xi = 1$ jet are relaxed compared to the $\xi = 2$ jet. The breaks in Figure 4 are due to the changing dependency of the jet bulk Lorentz factor on radius before and after saturation.

We have made several simplifying assumptions that are conservative in nature. First, we have adopted a fixed particle density in the entire the cocoon–jet boundary equal to the jet particle density. In reality, the density in the cocoon–jet boundary will be larger due to a particle density gradient from the more dense non-relativistic cocoon to the more tenuous jet. By fixing to the jet density we have been conservative in estimating the upper limit on $d\beta/dx$. We also neglected
$e^\pm$ pairs which at small radii will dominate over the $e^-$ associated with jet ions. As we detail in Section 3.2, the production of $e^\pm$ pairs aids in nuclei survival because the pairs provide additional targets to which nuclei lose energy, helping nuclei to cool on a shorter timescale. This would allow steeper velocity gradients. Finally, we have adopted the criterion that the spallation rate is smaller. This is perhaps conservative in light of the fact that a single spallation can still retain intermediate-mass nuclei composition. An alternative comparison we could have made would be to the spallation energy loss rate. Some authors have adopted other criteria, e.g., Metzger et al. (2011), and consider an optical depth $\lesssim 10$ to be acceptable. These relaxations will allow larger velocity gradients.

We have also neglected magnetic fields at the jet–cocoon boundary. Magnetic fields remain highly uncertain but may relaxations will allow larger velocity gradients.

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3. EFFECTS OF BULK ACCELERATION AND DISSIPATION

In the previous section, several candidates of the origin of jet nuclei were identified. However, even if nuclei are successfully loaded into the jet, they may be destroyed during the evolution of the jet because of bulk acceleration and/or dissipation. Here we discuss some of these effects.

3.1. Neutrons during Bulk Acceleration

Initially, the neutron component of the jet is well coupled to the ion component by elastic collisions with a small relative velocity $\beta \ll 1$ that is insufficient to cause nuclei to break up. However, the neutrons will lag behind the ions in the plasma during the bulk acceleration phase, if the collisions cannot keep up with the jet expansion. The velocity lag of neutrons relative to charged ions is $\Delta \Gamma/\Gamma \sim \tau_{\text{coll}}/\tau_{\text{exp}}$, where $\Gamma_n = \Gamma - \Delta \Gamma$ is the neutron Lorentz factor, $\tau_{\text{coll}} \approx 1/(n_i \sigma_{\text{ii}} \beta c)$ is the comoving collision timescale, $\sigma_i \approx \sigma_0 \bar{A}^{2/3}/\beta$ with $\sigma_0 \approx 3 \times 10^{-26}$ cm$^2$, and $\tau_{\text{exp}} \approx r/(\Gamma_n c)$ is the comoving expansion timescale.

The neutron–ion relative velocity is

$$\tilde{\beta} = \frac{\beta - \beta_n}{1 - \beta \beta_n} \sim \frac{\Gamma - \Gamma_n}{\Gamma} \sim \frac{\Gamma}{n_i \sigma_i \bar{A}^{2/3} \Gamma_n},$$

(26)

where $\beta_n$ is the neutron velocity. The relative velocity increases with radius during the bulk acceleration phase as $\tilde{\beta} \propto L_{\text{ke}}^{-1} r^{2-1} \eta$. Therefore, we require that at the end of the bulk acceleration phase ($r = r_{\text{sat}}$), the relative velocity is smaller than $\beta_{\text{sp}}$. This yields the condition for no spallation of

$$\eta < 220 L_{\text{ke}, 50}^{1/4} r_{0.7}^{-1/4},$$

(27)

for $\xi = 2$. To be conservative, we have made an Fe composition for the jet: this leads to a larger $\tilde{\beta}$ than a proton jet because of the longer collisional timescale. For a $\xi = 1$ jet, the saturation radius is larger by a factor $\eta$ and the expansion timescale is longer. Also, the higher jet density results in a faster collision timescale. Both serve to relax the nuclei survival condition, yielding $\eta < 1300 L_{\text{ke}, 50}^{1/3} r_{0.7}^{-1/3}$. Similar conclusions hold for our adopted magnetically dominated jet model: the survival condition is $\eta < 480 L_{\text{ke}, 50}^{5/3} r_{0.7}^{-2/3}$. If the above conditions are not satisfied, the neutron density is typically high enough that nuclei are spalled very rapidly (Beloborodov 2003). We plot the constraints for a $\xi = 2$ jet in Figure 5.

An alternative way for nuclei survival is if neutrons decouple from the accelerating plasma. We define the decoupling radius
We conclude that decoupling may play a role but its importance is increasing with acceleration, including spallation caused by collisions with neutrons (lines decreasing with $\eta$), for several values of $r_0$. Shaded region results in nuclei spallation (shade shown for our nominal $r_0 = 10^7$ cm and $\Delta T = 1$ s). (A color version of this figure is available in the online journal.)

As expected, decoupling becomes important for low-luminosity (LL) jets. For example, a small $L_{ke}$ strongly limits the range of $\eta$ for which spallation does not occur (Equation (27)), but the consideration of neutron decoupling allows a much wider range to be acceptable (Equation (31)). However, we caution the strong dependencies of Equation (31) on GRB parameters. We conclude that decoupling may play a role but its importance is strongly parameter dependent. Note that for a $\xi = 1$ jet the particle density is larger and the relative velocity $\beta$ is significantly slower, so that decoupling typically does not occur while the jet is propagating through the stellar core. In Figure 5, we show constraints on $L_{ke}$ and $\eta$ for several values of $r_0$, for a $\xi = 2$ jet with $\Delta T = 1$ s.

In conclusion, nuclei can survive over a significant range of $L_{ke}$ and $\eta$ because of slow neutron–ion relative velocities. For low $L_{ke}$ (and $\Delta T$), decoupling may occur before spallation so that a wider range of $\eta$ is allowed. For canonical values of $L_{ke}$ and $\eta$, collisions just reach spallation thresholds. For more collimated jets, a larger GRB parameter range is allowed.

3.2. Dissipation and Particle Heating at Oblique Shocks

Until now we have parameterized the jet cross section by Equation (4) which can be adapted to any morphology, e.g., conical (constant jet opening angle) or funnel (jet opening angle $\propto r^{-1/2}$) behavior. However, recent numerical and analytic works of jet propagation show that the jet morphology is not a free parameter but correlated with the jet’s interactions with the cocoon. According to the analytic work of Bromberg et al. (2011), which is based on the jet–cocoon solutions of Begelman & Cioffi (1989) and Matzner (2003), the jet morphology is determined by the jet energy density, its opening angle, and the density of the external medium. For a jet propagating through a dense material such as the progenitor star, the cocoon pressure is sufficiently strong and compresses the jet. Oblique collimation shocks form which generate the pressure needed to counterbalance the cocoon’s pressure. The shocks can form early in jet propagation but converge as the entire jet is shocked at a radius of approximately $r_{sh} \approx 0.1 r_h$, where $r_h$ is the jet head radius (e.g., Figure 2 of Bromberg et al. 2011).

For a sufficiently fast jet ($\Gamma_j > 1/\theta_j$), the Lorentz factor of the jet after being shocked is $\Gamma_j < 1/\theta_j$ (Bromberg et al. 2011), which for our GRB parameters $\theta_j \approx 0.13 \Omega_{de,-1}^{-1/2}$ and thus $\Gamma_j < 0.3 \Omega_{de,-1}^{-1/2}$. The jet ions passing the shock may obtain high random Lorentz factors corresponding to the relative Lorentz factor $\Gamma_j/(2\Gamma_r)$ in a similar way to the reverse shock (Section 2.3.2). Such ions would collide with other jet ions and any nuclei could be spalled. Here, we consider the spallation rate and energy loss rate to see whether jet nuclei are actually destroyed. To place conservative limits, we consider the fate of Fe ions within a jet plasma consisting of mainly protons. This assumption yields the highest spallation target density; the spallation rate is then $\nu_{sp} \approx n_{p} \sigma_{sp} c$, yielding $\nu_{sp}^{(2)} \sim 3 \times 10^{7} L_{ke,50,r_{sh},9^{-3/2}} n_{e}^{-1} \mathrm{s}^{-1} (\xi = 2)$ or $\nu_{sp}^{(1)} \sim 3 \times 10^{10} L_{ke,50,r_{sh},9}^{-3/2} \eta_{2}^{-1} \mathrm{s}^{-1} (\xi = 1)$.

To estimate the Fe–electron collisional cooling rate, we consider in addition to the electrons associated with jet protons the production of $e^\pm$ pairs. This directly increases the Fe cooling rate since $v_{ke} \propto n_{e}$. In fact, it is important for the survival of relativistic nuclei, since we showed in Section 2.3.3 that only mildly relativistic Fe ions can survive when electrons associated with jet protons are considered. For a radiation temperature $T_r$, the equilibrium pair density is (Shemi & Piran 1990),

$$n_{\pm} \approx 4.4 \times 10^{30} (T_r / m_e)^{3/2} e^{-m_e / T_r}.$$  (32)

We estimate the pair density in the jet by substituting the jet radiation temperature. For example, for a $\xi = 2$ jet in the bulk acceleration phase, $T_r^{(2)} \sim 180 L_{ke,50,r_{sh},9}^{-1} \mathrm{keV}$ and the resulting pair density exceeds those associated with jet protons by a factor of approximately $10^4$ at $10^8$ cm. However, by $\sim 10^{1.5}$ cm the $T_r$ has fallen too low and pairs make a negligible contribution to the total electron density. For a narrower jet,
The jet temperature remains high longer so that pairs make meaningful contributions out to larger radii, e.g., for $\xi = 1$ and canonical jet parameters, out to $\sim 10^{10}$ cm. Note that the jet radiation temperature can be high but photodisintegration is still slower than the spallation rate.

Let us first discuss when the fast jet approximation of Bromberg et al. (2011) is valid, i.e., $r_{sh}$ is larger than $8 \times 10^7$ cm ($\xi = 2$) or $6.4 \times 10^8$ cm ($\xi = 1$), both for our canonical choices of $r_0$ and $\Omega_{\psi}$. Whether collisional cooling prevents spallation depends quantitatively on the Lorentz factor of the Fe ions and the electron density. For example, at $r_{sh} = 10^9$ cm, the Fe ion cooling rate is $\nu_{s}\approx 1 \times 10^{12} E_{Fe, 12}^{-1} s^{-1}$ for our canonical GRB parameters and $\xi = 2$. Here we take the electron temperature to be equivalent to the radiation temperature because electrons reach local thermodynamic equilibrium on much shorter timescales compared to the ion–electron energy loss timescale. Comparing $\nu_s$ to the spallation rate $\nu_{sp}$, we see that Fe ions will undergo spallation if their energies are above $E_{Fe} \approx 4 \times 10^{13}$ eV, or a Lorentz factor of $\sim 800$. Thus, for our canonical parameters, Fe ions easily survive at $r_{sh} = 10^9$ cm.

However, as $r_{sh}$ increases, the pair density drops and survival becomes increasingly difficult. In Figure 6, we show the jet luminosity required for nuclei survival as a function of $r_{sh}$. This is determined by calculating the Fe Lorentz factor as a function of $r_{sh}$, and then the necessary pair density so that the energy loss rate $\nu_{s}$ is faster than the spallation rate $\nu_{sp}$. It is clear that higher luminosities are required for survival at larger radii, because the temperatures are lower. For a similar reason, a narrower jet is more favorable. A larger value of $\eta$ is only slightly more favorable mainly due to the smaller spallation target density it implies. Note that the importance of pairs can be noticed; for $L_{rad} = 10^{50}$ erg s$^{-1}$, nuclei survive out to $r_{sh} \approx 10^{8.5}$ cm ($\xi = 2$) or $r_{sh} \approx 10^{10}$ cm ($\xi = 1$); these radii are the radii where pairs cease to make a contribution. At the small radius end, the required luminosity suddenly drops because the Fe Lorentz factor tends to unity as the radius where the fast jet approximation becomes invalid is reached.

Similarly, jets with $\eta < 1/\theta_j$ do not satisfy the fast jet approximation, but oblique shocks are nonetheless expected as the jet expands into the stellar material and later on into the cocoon material. In this case we treat $\Gamma_s$ as a free parameter. If $\Gamma_s \sim \eta$, then nuclei are not spalled because the relative Lorentz factor between the jet and post-shocked jet is small. On the other hand, if $\Gamma_s \ll \eta$, the nuclei can attain mildly relativistic energies and can be spalled. For example, nuclei will survive in a $\eta = 5$ jet if $\Gamma_s \approx 2.5$ or higher. In Figure 6 we show as an example the case where $\Gamma_s$ is fixed such that nuclei are spalled (shown for $\Gamma_s = 2$; dotted). As expected, the smaller relative Lorentz factor results in a weaker requirement than our canonical ($\eta = 100$) jet of the same morphology (dashed).

Oblique shocks can also occur after jet collimation as the jet propagates through the progenitor, but the survival of nuclei depends largely on the post-shocked jet Lorentz factor. If $\Gamma_s \sim \eta$ as is usually expected, nuclei are likely not destroyed.

Note that in all of the above estimates we have conservatively assumed a proton-dominated jet and considered the survival of a (minor) Fe component. If the jet is mostly Fe dominated, the spallation target density decreases by $A = 56$ and improves Fe survival prospects. Also, we have conservatively considered the spallation rate. If we relax this and consider instead the spallation energy loss rate, i.e., we allow more than one spallation event, the limits for nuclei survival will be relaxed.

To conclude, jet nuclei may be spalled if jet ions obtain relativistic random velocities at oblique collimation shocks. However, when the shock radius $r_{sh}$ is small, the pair density in the jet can be sufficiently high that nuclei energy loss is faster than spallation. For our canonical jet parameters, this condition is realized for shock radii less than $\sim 10^{10}$ cm for an initially narrow jet (which may be realized if the jet is initially confined by the rotational funnel geometry of the progenitor) or less than $\sim 10^{10}$ cm for an initially conical jet. Nuclei that are entrained at larger radii (i.e., $r > r_{sh}$) may not be affected by the oblique collimation shocks and are not spalled. Jets with small $\eta$ imply small Fe ion Lorentz factors and are thus more favorable for survival.

Figure 6. Constraints on jet luminosity for Fe nuclei survival in collimation shocks. The shaded region results in nuclei spallation (shown for $a = 2$ and $\eta = 100$ jet). A narrower jet is more favorable for nuclei survival due to its higher temperature (dashed line). The $\eta = 5$ jet (dotted) is simply an illustration where the post-shocked jet Lorentz factor is deliberately fixed so that nuclei are spalled; it is possible that the post-shocked jet is faster in which case nuclei will not be spalled. For all other jets, the post-shock dynamics are calculated according to the collimation shock model of Bromberg et al. (2011); see the text.

3.3. Dissipation and Particle Acceleration at Emission Regions

GRB emissions typically consist of prompt gamma-ray emission and afterglow emission. The former is attributed to
non-thermal or quasi-thermal emissions produced via some internal dissipation, e.g., internal shocks or magnetic reconnections. In the standard optically thin synchrotron scenario of the internal shock model, inhomogeneities in the jet cause internal shocks within the jet, where the relative bulk kinetic energy is dissipated and accelerated electrons radiate gamma rays. In this scenario, it is natural to expect that ions are accelerated as well. The maximum energy to which ions are accelerated is roughly determined by requiring that the acceleration timescale is shorter than the dynamical and any energy loss timescales. Previous works have shown that the acceleration of nuclei to UHECR energies is possible for certain ranges of GRB parameters when internal dissipation happens in the optically thin regime (Murase et al. 2008; Wang et al. 2008).

The conditions required for the accelerated nuclei to survive in the prompt GRB photon field are much more stringent than those for acceleration. The comoving photodisintegration timescale for a nucleus moving through an isotropic photon background is (e.g., Murase et al. 2008; Wang et al. 2008)

$$\tau_{\text{photodisint}} = \frac{c}{2 \gamma A} \int_{\epsilon_b}^{\epsilon} d\epsilon \frac{\sigma_{\text{photodisint}}(\epsilon)}{\epsilon^2} \int_{2\gamma A}^{\infty} d\epsilon' \frac{1}{\epsilon'^2} \frac{dn}{d\epsilon},$$  

(33)

where quantities such as $\bar{\epsilon}$ are defined in the nucleus rest frame, $\gamma A$ is the nuclear Lorentz factor, $\epsilon_b \approx 7.6$ MeV is the threshold photon energy for photodisintegration of an iron nucleus, and $dn/d\epsilon$ is the differential photon spectrum. Adopting a comoving broken power law for the prompt GRB photon spectrum,

$$\frac{dn}{d\epsilon} \approx \epsilon^{-3/2} \epsilon_b^{-1} \left\{ \frac{\epsilon}{\epsilon_b} \right\}^{-1} < \epsilon < \epsilon_b < \epsilon < \epsilon_{\text{max}},$$  

(34)

where $\epsilon_b = \epsilon_{b,\text{obs}}/\eta$ is the break energy, the factor 5 in the denominator arises because the luminosity at the break energy is approximately one-fifth of the total luminosity, $\epsilon_{b,\text{obs}} \approx 1$ MeV, $\epsilon_{\text{min}} = 1$ eV, and $\epsilon_{\text{max}} = 10$ MeV, the comoving photodisintegration timescale is estimated to be

$$\tau_{\text{photodisint}} \approx 300 \epsilon_{b,\text{obs}}^{-1} L_{\text{ke}}^{-1/2} \eta \epsilon_{b,\text{obs}}^{-1} \text{s},$$  

(35)

where we have adopted $\xi = 2$ since the jet becomes conical after jet break out. We require that $\tau_{\text{photodisint}} < T_{\text{ion}}$ for UHECR nuclei to survive photodisintegration in the system, which yields the constraint

$$\eta > 110 \epsilon_{b,\text{obs}}^{-1/2} L_{\text{ke}}^{-1/2} \epsilon_{b,\text{obs}}^{-1} \text{s}^{-1}.$$

We plot this in Figure 7 as a function of the dissipation radius, i.e., emission radius. We emphasize that the optical depth for photodisintegration highly depends on the dissipation radius that is estimated to be $r_{\text{fs}} \approx 3c \delta t \epsilon^2 \eta^2$ for the internal shock case. Typical internal shock radii are $\sim 10^{30} - 10^{33}$ cm (e.g., Nakar & Piran 2002), so survival of UHECR nuclei is possible only for relatively large dissipation radii, although this is relaxed for sub-luminous GRBs (Murase et al. 2008). Typical radii for magnetic dissipation scenarios similarly span a wide range but can reach up to $\sim 10^{16}$ cm (e.g., Drenkhahn 2002; Granot et al. 2011; Metzger et al. 2011). In Figure 7 we show these ranges of dissipation radii for reference, noting that there are substantial uncertainties. Note that in the photospheric emission scenario, the emission radius is typically expected to be much smaller, where high-energy nuclei break up via both spallation and photodisintegration and instead high-energy neutrinos are generated (Murase 2008).

Dissipation and emission also occur later on, when the jet expands into the interstellar medium and sweeps up sufficient baryons to enter the Blandford–McKee phase. The forward and reverse shocks are collectively referred to as external shocks, where ions and electrons can be accelerated. The observed non-thermal photons are attributed to synchrotron emission from relativistic electrons accelerated at such shocks. Compared to the case of the prompt emission, nuclei interact with softer photons, and it has been shown that the survival of UHECR nuclei is also possible (Murase et al. 2008; Wang et al. 2008).

4. DISCUSSIONS AND CONCLUSIONS

In this work, we focus on the fate of nuclei in a relativistic jet accompanying a CCSN. We consider three sources of jet nuclei (1) loading at the jet base during jet launch, (2) in situ explosive nucleosynthesis, and (3) entrainment of external nuclei during jet propagation. We first discuss the conditions for nuclei survival in each of them. For initial loading, we find that nuclei survive if the jet launch radius is greater than $r_0 \sim 10^8$ cm or the initial radiation luminosity is less than $L_{\text{rad}} \sim 10^{48}$ erg s$^{-1}$ (Figure 2). These conditions may be satisfied by magnetic models of classic GRBs or models of sub-luminous GRBs. If nuclei are destroyed at jet launch into free nucleons, the material is fused into nuclei, but the freezeout abundance contains significant heavy nuclei only for low entropy jets (Beloborodov 2003; Inoue et al. 2003; Metzger et al. 2011). Finally, whether external nuclei survive during entrainment into the jet depends critically on the velocity gradient at the cocoon–jet boundary (Figure 4). If the gradient is too steep, collisions between external nuclei and jet ions become too energetic and external nuclei are spalled. On the other hand, if the gradient is shallow enough, the nuclei thermalize with the jet plasma before spallation and survive. Since the growth rate of shear-driven instabilities at the cocoon–jet boundary is larger than the nuclei energy loss or spallation times, the velocity gradient is expected to be shallow enough to allow nuclei to survive.
We find that it is strongly parameter dependent but tends to be model-dependent. Typical parameters are taken to be $L_{\text{ke}} = 10^{50}$ erg s$^{-1}$, $\eta = 300$, $\xi = 2$, and $r_0 = 10^7$ cm for HL GRB and $L_{\text{ke}} = 10^{47}$ erg s$^{-1}$, $\eta = 10$, $\xi = 2$, and $r_0 = 10^7$ cm for LL GRB. For entrainment, we consider the currently unconstrained jet–cocoon velocity gradient to be model-dependent.

We emphasize that we have attempted to place conservative constraints on nuclei survival. This means we have not always used the same jet composition in each investigation. For example, when we consider neutron collisions we adopt a nuclei-dominated jet so that the relative velocity is highest and spallation occurs most readily; and when we consider oblique shocks we adopt a proton-dominated jet so that the spallation rate is maximal. In the same spirit, we assume that a single spallation process can still be effective. The low $\eta$ also works positively for nuclei survival in oblique collimation shocks, because the nuclei Lorentz factors are small. However, note that nuclei can still be spalled at large $r_{sb}$ if $\Gamma_s \ll \eta$. Finally, though uncertainty is large, survival at dissipation is easier than in LH GRBs (Figure 7; see also Murase et al. 2008; Wang et al. 2008).

As seen in this work, jets accompanying CCSN may contain heavy nuclei that originate from the stellar core, disk wind, and/or low entropy (kinetically or magnetically dominated) jets. A high abundance of nuclei is attractive in view of recent reports of a heavy-ion-dominated composition of the highest-energy UHECR. If AGNs are the sources of UHECRs, rigidity-dependent acceleration models naturally predict a heavy-ion composition at the highest energies, but this seems to be inconsistent with the null observation of an excess of protons at lower energies (Abreu et al. 2011b). On the other hand, nuclei-rich UHECRs could be realized in GRBs and CCSN, as long as the accelerated particles are injected without being broken up.

Next we investigate the conditions for nuclei survival once loaded into the jet. First, collisions with neutrons become energetic during the jet bulk acceleration phase and can cause nuclei spallation if the GRB terminal Lorentz factor is larger than $\eta \sim 220 L_{\text{ke,50}}^{1/4} r_{0,7}^{-1/4}$ (for a conical jet) or $\eta \sim 1300 L_{\text{ke,50}}^{1/3} r_{0,7}^{-1/3}$ (for a funnel jet); for smaller $\eta$, spallation is energetically prohibited. Also, we identify the parameter space where the neutrons decouple before causing spallation. We find that it is strongly parameter dependent but tends to help with small $L_{\text{ke}}$ jets (Figure 5). Second, we show that nuclei can survive at oblique shocks if $e^{\pm}$ pairs are produced such that the nuclei collisional energy loss rate is competitively high. Such conditions are typically realized for shock radii $r_{sb}$ less than $\sim 10^8-10^{10}$ cm (Figure 6). A smaller $\eta$ works positively for nuclei survival, as does a post-shocked jet Lorentz factor $\Gamma_s$ that is close to the pre-shocked jet Lorentz factor.

Based on the above results, we can consider the sources and survival of nuclei in three distinct sets of GRB models. These are summarized below and in Table 1.

First, we conclude that nuclei from initial loading and explosive nucleosynthesis are disfavored in the HL GRB fireball scenario (see also Beloborodov 2003); instead, nuclei may come from entrainment during jet propagation. Once loaded, nuclei may survive collisions with neutrons. Survival in oblique collimation shocks is possible while $r_{sb}$ is less than $\sim 10^7-10^9$ cm. After $r, p, m, \eta_{sb}$ reaches these values, nuclei entrained at smaller radii ($r < r_{sb}$) may be spalled as they cross the shocks, while nuclei that are entrained at larger radii ($r > r_{sb}$) may survive. Successfully entrained nuclei can survive the emission region provided the dissipation radius is large enough (Figure 7; see also Murase et al. 2008; Wang et al. 2008).

On the other hand, multiple sources of nuclei are possible for magnetic-dominated scenarios of HL GRBs. For example, nuclei may survive at initial loading, because for a given GRB luminosity the initial radiation energy can be lower. Even if nuclei are spalled at loading, the subsequent explosive nucleosynthesis can lead to Fe-group elements and possibly beyond (Metzger et al. 2011). Entrainment can work too, if, as in recent models, the jet is collimated by a cocoon. The slower bulk acceleration compared to fireballs means collisions with neutrons are generally less destructive, and oblique shock are expected to be weaker, although definite statements are dependent on the specific model of magnetic dissipation.
that loading and entrainment of external baryons occur efficiently. Jets accompanying CCSN can have a wide range of baryonic content, with GRB jets among the cleanest with fewest baryons. GRB jet models achieve this in different ways. For example, magnetic-dominated jets that are powered by energy release on field lines where baryons are confined can plausibly contain very few initial baryons (e.g., Lyutikov & Blandford 2003). Thus, although nuclei are more likely to be present at the launch sites of magnetically dominated jets (Table 1), they may not be efficiently loaded. However, we stress that such a statement remains highly model dependent. Later entrainment is attractive in the sense that jets must load baryons at some point, and propagation through the baryon-rich progenitor is an unavoidable course of jet evolution. The entrainment efficiency is expected to depend on the instability mechanism, thermodynamic parameters of the cocoon and jet, magnetic field, and so on, and thus evolve. To make quantitative estimates of these and other effects requires a detailed model of entrainment and is beyond the scope of the present paper. In this work, we investigated the process by which nuclei may survive during loading. Although the specific GRB model may change in the future, the physical processes by which nuclei survive should still hold.

We can, however, speculate about the composition assuming entrainment is the main process of jet baryon loading. In an oversimplified picture where entrainment occurs equally efficiently during the entire propagation of the jet, the jet composition will be similar to the total matter swept up in the cocoon. Indeed, high $^{56}$Ni abundance of the swept up material. Although the specific GRB model may change in the future, the physical processes by which nuclei survive should still hold.

Finally, it is relevant to identify observational signatures of heavy nuclei. For protons, neutrons, and protons interact with photons and/or target protons and produce charged pions. The neutrino signatures have been well studied in various contexts. If collisionless shocks can be formed inside the progenitor star, GeV–TeV neutrinos might be expected during jet propagation (e.g., Mészáros & Waxman 2001; Razzanu et al. 2003, 2004; Ando & Beacom 2005; Horiiuchi & Ando 2008; Iocco et al. 2008). One may expect TeV–PeV neutrino emission associated with photospheric and/or shock breakout emissions, where the formation of collisionless shocks is expected (Murase 2008; Murase et al. 2011; Katz et al. 2011). At larger dissipation radii above the photosphere, PeV–TeV neutrinos might be produced when the oblique shock radius exceeds $R_{sh} \sim 10^{10}$ cm. Nuclei that lie at radii $R_{sh} \gtrsim 10^{10}$ cm, mainly intermediate nuclei such as carbon, are thus more likely to be spalled before entering the cocoon.

The existence of nuclei can also be probed by identifying atomic or nuclear line emissions. As in discussed Mészáros & Rees (2001), the cocoon material produced by the jet would break out of the stellar envelope, where iron-enriched clumps may be shed by a UV/X-ray continuum and an Fe line luminosity of $\sim 4 \times 10^{47}$ erg s$^{-1}$($M_\odot / 10^{-5} M_\odot$) is expected. Here $M_\odot$ is the bubble mass and $x_{fe}$ is the Fe mass fraction. Alternatively, the Fe line emission may be caused by X-rays produced by a continuous but decaying jet (Rees & Mészáros 2000). Though the details are uncertain, if the high abundance of heavy nuclei is realized, X-ray line features might be constrained or even detected by current and future X-ray satellites such as Astro-H. Furthermore, it is possible that boosted atomic photons could be expected in the GeV range for accelerated nuclei (Kusenko & Voloshin 2011).

Nuclei can also emit $\sim$ MeV gamma rays (in their rest frame) through their excitation states. Such nuclear gamma rays are not so easy to detect for distant GRBs, but they may be detected for nearby and/or luminous events. In particular, nuclear gamma rays may be boosted to the TeV energy range when nuclei are accelerated. For example, the radioactive isotopes $^{56}$Ni and $^{56}$Co lead to the production of MeV nuclear gamma-ray lines after they are excited via, e.g., electron capture or positron emission (for details, see, e.g., Milne et al. 2004; Horiiuchi & Beacom 2010). These nuclei can be entrained and lead to GeV–TeV gamma rays in GRBs and hypernovae (Ioka & Mészáros 2010). The timescale of emission is also extended by the nuclei Lorentz factor to $\sim 10^5 \gamma_{A,5}$ for $^{56}$Co (where $\gamma_{A}$ is the Lorentz factor of the nuclei in the observer frame), and detection is possible only for Galactic events. On the other hand, nuclei interact with low-energy photons in the source via the photodisintegration process, leading to subsequent (almost prompt) $\sim 0.2$ TeV $\gamma_{A,5}$ gamma rays from daughter excited nuclei. This signal is useful as a unique probe of nuclei acceleration as well as synchrotron and inverse-Compton gamma rays from pairs generated via the Bethe–Heitler process (Murase & Beacom 2010a; Aharonian & Taylor 2010). There are a multitude of potential signatures to be explored.

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