FINDING THE FIRST COSMIC EXPLOSIONS. II. CORE-COLLAPSE SUPERNOVAE

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ABSTRACT

Understanding the properties of Population III (Pop III) stars is prerequisite to elucidating the nature of primeval galaxies, the chemical enrichment and reionization of the early intergalactic medium, and the origin of supermassive black holes. While the primordial initial mass function (IMF) remains unknown, recent evidence from numerical simulations and stellar archaeology suggests that some Pop III stars may have had lower masses than previously thought, 15–50 $M_\odot$ in addition to 50–500 $M_\odot$. The detection of Pop III supernovae (SNe) by JWST, WFIRST, or the TMT could directly probe the primordial IMF for the first time. We present numerical simulations of 15–40 $M_\odot$ Pop III core-collapse SNe performed with the Los Alamos radiation hydrodynamics code RAGE. We find that they will be visible in the earliest galaxies out to $z \sim 10–15$, tracing their star formation rates and in some cases revealing their positions on the sky. Since the central engines of Pop III and solar-metallicity core-collapse SNe are quite similar, future detection of any Type II SNe by next-generation NIR instruments will in general be limited to this epoch.

Key words: early universe – galaxies: high-redshift – hydrodynamics – radiative transfer – shock waves – stars: early-type – supernovae: general

Online-only material: color figures

1. INTRODUCTION

The first stars are crucial to the formation of primeval galaxies (Johnson et al. 2008, 2009; Greif et al. 2008, 2010; Jeon et al. 2012; Pawlik et al. 2011, 2013; Wise et al. 2012), the chemical enrichment of the early intergalactic medium (IGM; Mackey et al. 2003; Smith & Sigurdsson 2007; Smith et al. 2009; Chiaki et al. 2013; Ritter et al. 2012), the initial stages of cosmological reionization (Whalen et al. 2004; Kitayama et al. 2004; Alvarez et al. 2006; Abel et al. 2007; Wise & Abel 2008), and the origin of supermassive black holes (Bromm & Loeb 2003; Johnson & Bromm 2007; Djorgovski et al. 2008; Milosavljević et al. 2009; Alvarez et al. 2009; Lippai et al. 2009; Tanaka & Haiman 2009; Li 2011; Park & Ricotti 2011, 2012a, 2012b; Johnson et al. 2012a, 2012b; Whalen & Fryer 2012; Agarwal et al. 2012). Unfortunately, because they lie beyond the reach of current ground- and space-based instruments there are not yet any observational constraints on their properties. The early numerical simulations of Population III (Pop III) stars suggest that they form in isolation, one per halo, in $10^5$–$10^6$ $M_\odot$ dark matter halos at $z \sim 20–30$ and are very massive, 100–500 $M_\odot$ (Bromm et al. 1999, 2002; Abel et al. 2000, 2002; Nakamura & Umemura 2001; Yoshida et al. 2008). Such stars usually drive out the baryons from their halos in strong ionized flows that create diffuse H II regions with $n \sim 0.1–1$ cm$^{-3}$ (Whalen et al. 2004; Kitayama et al. 2004; Alvarez et al. 2006; Abel et al. 2007; Wise & Abel 2008; Whalen & Norman 2008a, 2008b). Pop III stars are not thought to lose much mass over their lifetimes because there are no line-driven winds in their metal-free atmospheres (Kudritzki 2000; Vink et al. 2001; Baraffe et al. 2001; Krtička & Kubát 2006; Ekström et al. 2008).

There is growing evidence that some Pop III stars may be less massive than previously thought. Recent, more extensive ensembles of numerical simulations have found many halos with central collapse rates consistent with 20–60 $M_\odot$ for the final mass of the star (O’Shea & Norman 2007, 2008; Wise & Abel 2007) and that a fraction of the halos form binaries in this mass range (Turk et al. 2009). New simulations of Pop III protostellar accretion disks suggest that they were prone to fragmentation into as many as a dozen smaller stars (Stacy et al. 2010; Clark et al. 2011; Smith et al. 2011; Greif et al. 2011, 2012). Preliminary models also suggest that ionizing UV breakout from primordial protostellar disks limits the masses of some Pop III stars to ∼40 $M_\odot$ (Hosokawa et al. 2011, 2012; Stacy et al. 2012; although also see Omukai & Palla 2001, 2003; Omukai & Inutsuka 2002; Tan & McKee 2004; McKee & Tan 2008).

While these new models are important steps forward they cannot predict the final masses of primordial stars. For example, the fragmentation of Pop III protostellar disks has only been followed out to a few centuries, far short of the time required to build up a massive star, and many of the fragments in these models are later found to merge with the central object instead of becoming stars themselves. Thus, it is not clear if a small swarm of less massive stars forms or a single, very massive star is created, albeit through protracted clumpy accretion. Clumpy accretion may also keep the protostar puffed up and cool, allowing it to reach much higher masses before evaporating its accretion disk than in the Hosokawa et al. (2011) models. Because no simulation realistically bridges the gap between the formation and fragmentation of the accretion disk and its destruction up to a Myr later, they cannot yet constrain the Pop III initial mass function (IMF; for recent reviews, see Glover 2013; Whalen 2012). We also note that
the effects of turbulence (Latif et al. 2013), magnetic fields due to the small scale turbulent dynamo (Schober et al. 2012), and radiation transport on the evolution of the disk are not well understood. There have been attempts to determine the masses of the first stars by comparing the cumulative nucleosynthetic yield of their supernovae (SNe) to the fossil abundance record, the chemical abundances measured in ancient, dim extremely metal-poor (EMP) stars out in the Galactic halo (e.g., Beers & Christlieb 2005; Frebel et al. 2005; Caffau et al. 2012). Pop III stars from 15–40 $M_\odot$ die in core-collapse (CC) SN explosions and 140–260 $M_\odot$ stars explode as far as more energetic pair-instability (PI) SNe (Heger & Woosley 2002; although the lower mass limit for PI SNe has recently dropped down to 65 $M_\odot$ for rotating progenitors; Chatzopoulos & Wheeler 2012). Joggerst et al. (2010, hereafter JET10) found that the chemical yields of 15–40 $M_\odot$ Pop III SNe have now been found in high redshift damped Ly$\alpha$ absorbers (Cooke et al. 2011), and there are indications it could be present in a new sample of 18 metal-poor stars in the Sloan Digital Sky Survey (Ren et al. 2012; see also Karlsson et al. 2008 on why the odd–even effect may not have been discovered in earlier surveys). Together, these discoveries suggest that both high-mass and low-mass Pop III star formation was possible. In sum, while stellar archaeology has revealed some insights into the first stellar populations it has not yet placed firm constraints on their masses. Observations of primordial SNe (Bromm et al. 2003; Vasiliev et al. 2012) could directly probe the Pop III IMF for the first time because in principle they can be detected at great distances and distinguish between low mass and high mass progenitors. Until now, calculations of Pop III SN light curves and spectra have been confined to PI SNe to determine if they can be detected in high redshift damped Ly$\alpha$ absorbers (Cooke et al. 2011), and there are indications it could be present in a new sample of 18 metal-poor stars in the Sloan Digital Sky Survey (Ren et al. 2012; see also Karlsson et al. 2008 on why the odd–even effect may not have been discovered in earlier surveys). Together, these discoveries suggest that both high-mass and low-mass Pop III star formation was possible. In sum, while stellar archaeology has revealed some insights into the first stellar populations it has not yet placed firm constraints on their masses.

2. NUMERICAL ALGORITHM

We adopt the models in JET10 for our grid of light curve simulations because they span the range of progenitor masses, structures, and explosion energies expected for such stars and because their elemental yields are a good match to those in the fossil abundance record. We calculate light curves and spectra in four stages. First, 15–40 $M_\odot$ Pop III stars are evolved through all stages of stable nuclear burning and then exploded in the Kepler code. Each explosion is then mapped onto a two-dimensional (2D) adaptive mesh refinement (AMR) grid in the CASTRO code and evolved until just before the shock reaches the surface of the star to capture internal mixing prior to breakout. We then spherically average our CASTRO profiles onto one-dimensional (1D) AMR grids in RAGE and evolve them out to four months, after which the ejecta dims below observability. Finally, we post-process our RAGE profiles with the SPECTRUM code to construct light curves and spectra with detailed atomic opacity sets.

2.1. Kepler

We use the 1D Lagrangian stellar evolution code Kepler (Weaver et al. 1978; Woosley et al. 2002) to evolve the progenitors from the beginning of the main sequence to the onset of collapse of their iron cores. The explosions are then artificially triggered by a piston at a constant Lagrangian mass coordinate that advances through the star with a specified radial history. The SNe are evolved through all nuclear burning ($\sim$100 s after the explosion) and then the simulations are halted, well before the shock exits the He shell and any reverse shocks form in which instabilities might develop. We follow energy production with a 19 isotope network until oxygen is depleted from the core and with a 125 isotope quasi-equilibrium network thereafter. Stellar rotation is approximated by a semi-convective mixing parameter.

As in JET10, we consider two progenitor metallicities ($Z = 0$ and $10^{-4} Z_\odot$, denoted by $z$ and $u$, respectively), three explosion energies ($E_{\text{ex}} = 0.6$, 1.2, and 2.4 $B$ designated by $B$, $D$, and $G$, respectively, where $1 B = 1$ Bethe = 10$^{53}$ erg), and three progenitor masses (15, 25, and 40 $M_\odot$), a total of 18 models. Thus, model z15G is the 2.4 $B$ explosion of a zero-metallicity $15 M_\odot$ Pop III star. We use only one rotation rate from that study ($R = 5$, where $R$ is defined in Section 3 of JET10) because it was found that the rate of rotation had little effect on the degree of mixing during the explosion. In contrast to the u- and z-series 150–250 $M_\odot$ stars in Whalen et al. (2012a), u-series 15–40 $M_\odot$ Pop III stars die as compact blue giants and z-series stars die as red supergiants.

Further details on our Kepler models, such as our numerical mesh, explosion triggers, and use of semi-convective mixing as a proxy for rotation, can be found in Section 2 of JET10. Besides altering the envelope of the star, rotation can also trigger asymmetries in the central explosion engine whose effects on the SN light curves are not explored here. For example, jet-like instabilities could arise in the shock that lead to breakout times that vary with angle and alter the luminosity of the explosion. In extreme cases, such jets can burn material in the core and form large amounts of $^{56}$Ni that also cause the SN to be brighter at intermediate times (e.g., hypernovae; Iwamoto et al. 2005). Asymmetries in the explosion can be seeded prior to core bounce and cause even stronger instabilities and central mixing at later times. The effects of these processes will be examined in future models.
describe the physics implemented in our RAGE models in Frey et al. (2013, hereafter FET12): multispecies advection, 2-temperature (2T) gray flux-limited diffusion, and energy deposition due to the radioactive decay of $^{56}$Ni (Fryer et al. 2009). We note that 2T radiation transport, in which radiation and matter temperatures are evolved separately, better models shock breakout and its aftermath than previous 1T models of SN explosions. We exclude gravity from our RAGE calculations because JET10 showed that fallback onto the central remnant is over before breakout and that its gravity does not strongly affect the shock after it reaches the surface of the star. We evolve mass fractions for the same 27 elements as in CASTRO.

2.3.1. Model Setup

We spherically average 2D densities, velocities, temperatures and mass fractions for the explosion, the star, and the circumstellar envelope from CASTRO onto a 50,000 zone 1D spherical AMR mesh in RAGE. Since we do not evolve radiation energy densities (erg cm$^{-3}$) in CASTRO, we initialize them in RAGE by assuming that

$$e_{\text{rad}} = aT^4,$$

where $a$ is the radiation constant and $T$ is the gas temperature in CASTRO. Also, we construct the specific internal energy (erg g$^{-1}$) from $T$ from

$$e_{\text{gas}} = C_V T,$$

where $C_V = 7.919 \times 10^7$ erg K$^{-1}$ is the specific heat of the gas. This prescription does not fully hold in the cool outermost layers of red supergiants where the temperature falls below $\sim$1 eV because it excludes ionization energies but is adopted there for simplicity. The initial radius of the shock varies from $4.1 \times 10^{13}$ to $1.3 \times 10^{14}$ cm in the $z$-series and $2.04 \times 10^{12}$ to $4.77 \times 10^{12}$ cm in the $u$-series. Our root grid at setup has a resolution of $8.0 \times 10^6$ cm and an outer boundary at $4.0 \times 10^{14}$ cm for the $z$-series and a resolution of $1.0 \times 10^9$ cm and an outer boundary at $5.0 \times 10^{15}$ cm in the $u$-series. We allow up to five levels of refinement in both the initial interpolation of the profiles onto the setup grid and later throughout the simulation. Our setup guarantees that all important features in the profiles are resolved by at least 10 grid points and can be subsampled by up to 32 times more points if necessary. We ensure that the explosion is spanned by at least 5000 zones at shock breakout so that the photosphere of the shock is fully resolved; failure to do so can lead to underestimates of luminosity during post-processing.

We impose reflecting and outflow boundary conditions on the fluid and radiation flows at the inner and outer boundaries of the mesh, respectively. When the calculation is launched, Courant times are initially short due to high temperatures, large velocities, and small cell sizes. To speed up the simulation and to accommodate the expansion of the ejecta, we periodically resample the profiles onto a larger mesh as the explosion grows. Each regrid significantly increases the time step on which the model evolves. We remap just the explosion itself, excluding any medium beyond the shock, and then graft the original environment that lies beyond this radius onto the shock on the new grid. We apply the same criteria in choosing a new root grid as in the original problem setup: any important density or velocity structure must be resolved by at least 10 mesh zones and be sampled by up to 32 times more zones as needed, and at least 5000 coarse grid zones are allocated to the blast after breakout. The outer boundary of the final largest mesh in our

2.2. CASTRO

As in JET10, we map our Kepler explosion profiles onto a 2D axisymmetric grid in the new CASTRO code (Almgren et al. 2010) and evolve the shock up to the edge of the star in mass coordinate. CASTRO is a multidimensional Eulerian AMR code with a high order, unsplit Godunov hydrodynamics scheme. We follow the same 27 elements from hydrogen to zinc as in JET10 and include the gravity of the compact remnant, which JET10 and others have shown to be crucial to both fallback and mixing in Pop III CC explosions (Zhang et al. 2008). Radiation transport is not necessary in this stage of our calculations because photon mean free paths in the star prior to breakout are so short that they are simply advected through the star by the fluid flow. We do account for the contribution of photons to pressure in the equation of state (EOS). Our models also include energy deposition due to radioactive decay of $^{56}$Ni in the ejecta, as described by Equation (4) in JET10.

The JET10 models show that mixing is mostly complete by the time the shock reaches the surface of the star, so elemental mass fractions are realistically distributed in radius and angle in our profiles when we spherically average them and map them into RAGE. We halt the CASTRO simulation when the shock is no less than 100 photon mean free paths $\lambda_p$ from the edge of the star:

$$\lambda_p = \frac{1}{\kappa \rho},$$

where $\kappa$ is the opacity due to Thomson scattering from electrons, taken to be 0.285 g cm$^{-2}$, and $\rho$ is the density just ahead of the shock inside the star.

This intermediate step in CASTRO allows us to capture how heavy elements that have built up in the star are mixed in the initial stages of the SN without resorting to full 2D radiation transport models in RAGE. Such calculations, while tractable, would require too much time for a grid of models as large as ours and are not necessary because photon transport does not affect mixing. Mixing determines the order in which emission and absorption lines appear in the spectra over time. If metal lines appear soon after shock breakout in a Pop III SN they would be absorption lines appear in the spectra over time. If metal lines appear soon after shock breakout in a Pop III SN they would be absorption lines appear in the spectra over time.
models is $5.0 \times 10^{16}$ cm, the greatest distance the ejecta can propagate in four months.

### 2.3.2. Circumstellar Envelope

The ambient media of $z \sim 30$ and $z \sim 10–15$ Pop III PI SNe are quite different because the former occur in small cosmological halos and the latter go off in primeval galaxies. Even low-mass Pop III stars can photoevaporate minihalos (Whalen et al. 2008c), and so at $z \gtrsim 20$ most die in diffuse relic H II regions with $n \sim 0.1$ cm$^{-3}$. Local densities in $z \sim 15$ protogalaxies are less well understood but are likely higher. In either case, if the star sheds a wind it will determine the density profile closest to the star. When considering PI SNe, Whalen et al. (2012a) allowed for the possibility that massive Pop III progenitors have modest winds in addition to being surrounded by a diffuse H II region. However, 15–40 $M_\odot$ Pop III stars are less likely to drive strong winds given their lower surface temperatures, luminosities, and the fact that stars in this mass range today manifest only weak winds, even at solar metallicities. Consequently, in this study we join a very low-mass wind profile to the surface of the star:

$$\rho_W(r) = \frac{\dot{m}}{4\pi r^2 v_W},$$

where $\dot{m}$ is the mass loss rate of the wind and $v_W$ is its speed.

### 4. SPECTRUM

To calculate a spectrum from a RAGE profile, we first map its densities, temperatures, velocities, and mass fractions onto a 2D grid in the Los Alamos SPECTRUM code. SPECTRUM performs a direct sum of the luminosity of every fluid element in this discretized profile to compute the total flux emitted by the ejecta along the line of sight at every wavelength. This procedure is described in FET12, and it accounts for Doppler shifts and time dilation due to the relativistic expansion of the ejecta. SPECTRUM also calculates the intensities of emission lines and the attenuation of flux along the line of sight, thereby capturing both limb darkening and absorption lines imprinted on the flux by intervening material in the ejecta and wind.

Gas densities, velocities, mass fractions and radiation temperatures from the finest levels of refinement on the RAGE AMR grid are first extracted and ordered by radius into separate files, with one variable in each file. Because of limitations on machine memory and time, we map only a subset of these points onto the SPECTRUM grid. We sample the RAGE radiation energy density profile inward from the outer boundary to determine the position of the radiation front, where $aT^4$ rises above 1.0 erg cm$^{-3}$. The maximum depth from which radiation can escape the ejecta is then determined. For radiating shocks this optical depth is often taken to be

$$\tau = \frac{c}{v_{sh}},$$
where \(c\) and \(v_{sh}\) are the speed of light and the shock, respectively. This limit is derived from the theory of radiating shocks by assuming that the diffusion limit holds for photons in the shock and defining their diffusion timescale through the shock to be (e.g., Ohyama 1963)

\[
\tau_{\text{diff}} = \left( \frac{\Delta x}{\lambda} \right)^2 \frac{\lambda}{c},
\]

(7)

where \(\Delta x\) is roughly the width of the shock and \(\lambda\) is the photon mean free path. Setting the diffusion velocity of the photons through the shock to be

\[
v_{\text{diff}} = \frac{\Delta x}{\tau_{\text{diff}}},
\]

(8)

and recalling that

\[
\tau = \frac{\Delta x}{\lambda},
\]

(9)

it follows that

\[
v_{\text{diff}} = \frac{c}{\tau}.
\]

(10)

If one equates the photon diffusion velocity with \(v_{sh}\), it follows that the greatest optical depth from which photons can escape the shock is given by Equation (6). This derivation also assumes that the opacity is constant. For the 50,000 km s\(^{-1}\) flows typical of SNe, \(\tau \sim 6\).

However, the assumption of the diffusion limit and constant opacity likely does not fully hold, particularly during shock breakout when photons are no longer trapped by the shock, so we adopt \(\tau = 20\) as a more careful limit. We find the radius of the \(\tau = 20\) surface in the ejecta by integrating the optical depth due to Thomson scattering inward from the outer boundary, taking \(\kappa_{\text{Th}}\) to be 0.288 for H and He gas at primordial composition.

The extracted gas densities, velocities, temperatures, and species mass fractions are then interpolated onto a 2D grid in \(r\) and \(\theta\) in SPECTRUM whose inner boundary is that of the RAGE mesh and whose outer boundary is 10\(^{18}\) cm. One hundred uniform zones in log radius are assigned from the center of the grid to the \(\tau = 20\) surface, and the region between the \(\tau = 20\) surface and the edge of the radiation front is partitioned into 5000 uniform zones in radius. The wind between the front and the outer edge of the grid is divided into one hundred uniform zones in log radius, for a total of 5200 radial bins. The data within each of these new radial bins is mass-averaged to guarantee that the SPECTRUM grid captures very sharp features from the original RAGE profile. The grid is uniformly divided into 160 bins in \(\mu = \cos \theta\) from \(-1\) to 1. Our mesh fully resolves regions of the flow from which photons can escape the ejecta and only lightly samples those from which they cannot.

The sum of the luminosities over all wavelengths in one spectrum yields the bolometric luminosity of the SN at that moment. Many such luminosities computed over a range of times constitutes the light curve of the explosion. We cover shock breakout with 50 spectra uniformly spaced in time and the rest of the light curve with 200 spectra that are logarithmically distributed in time out to four months. Each SPECTRUM calculation requires 3–6 hr on 32 processors on LANL platforms.

3. BLAST PROFILES, LIGHT CURVES, AND SPECTRA

Hydrodynamical profiles for shock breakout are shown for the fiducial cases u40G, a blue giant progenitor, and z40G, a red supergiant progenitor, in Figures 2 and 3. The reverse shock in the z40G velocity profile in Figure 1, which vigorously mixes the interior of the star prior to breakout (Figure 5 of Joggerst et al. 2010), is gone by the time the shock reaches the surface of the star at 3.91 \times 10^5 s. With no reverse shock there is no further Rayleigh–Taylor mixing, and the metals are essentially frozen in mass coordinate as the ejecta expands. Our procedure for spherically averaging 2D CASTRO mass fractions to approximate their radial distribution in 1D in RAGE therefore
Figure 3. Shock breakout in the z40G (red supergiant) run. Left panel: from left to right, velocities at $3.91 \times 10^5$ s, $3.97 \times 10^5$ s, and $4.14 \times 10^5$ s. Right panel: gas temperatures at the same times.

The breakout transient is visible as the flat plateau in gas temperature that moves ahead of the ejecta. This plateau traces the leading edge of the radiation front. The height of the plateau is the temperature to which the radiation pulse heats the gas as it passes through it, not the temperature of the fireball itself, which is much higher. The u40G shock is much hotter than the z40G shock at breakout ($\sim 100$ eV versus 10 eV) because it has broken out of a much more compact star and done less $PdV$ work on its surroundings. As a result, the radiation front initially heats the wind to 20 eV in u40G but to only 1 eV in z40G. As the shock expands, its spectrum softens, and the temperature of the gas behind the radiation front falls over time: from 20 eV to 8 eV in u40G from $1.81 \times 10^5$ s to $1.93 \times 10^5$ s and from 1.2 eV to 0.4 eV in z40G from $3.91 \times 10^5$ s to $4.14 \times 10^5$ s.

In both explosions, shock breakout coincides with radiation breakout because the moderate wind density at the surface of the star cannot confine the radiation front. As both shocks descend the $r^{-2}$ wind gradient, the u40G SN accelerates but the z40G SN speeds up and then decelerates. The u40G shock continues to accelerate even though ambient densities are 1000 times higher at the surface of the star because the explosion is much hotter and the shock is more radiatively driven than the much cooler z40G shock. Soon after breakout, both SNe assume self-similar free expansion velocity profiles and homologously expand, as we show in Figure 4 for u40G. This self-similar expansion continues until the end of the simulations at four months. Eventually, spherical dilution of the ejecta renders the SN transparent, and it enters the nebular phase. At this stage, fluid elements in the ejecta may no longer remain in radiation equilibrium with their neighbors so the true opacity may deviate from the OPLIB opacities, which assume local thermodynamic equilibrium (LTE). Also, the Kirchhoff–Planck relation (Equation (4) of Frey et al. 2013) may not fully hold in the SPECTRUM code. There may therefore be some inaccuracy in our spectra at late times, but these events are only visible in the NIR at much earlier stages of the explosion, as we later show.

We note that early simulations of shock breakout exhibited a spike in the matter temperature ahead of the shock (Ensmann & Burrows 1992). Our simulations assume that the ions and electrons are strongly coupled, leading to only mild deviations between the matter and radiation temperatures. As such, we do not expect, and our simulations do not produce, a strong spike in the matter temperature at the front of the shock, in agreement with many recent results (Tominaga et al. 2011; Moriya et al. 2013; Tolstov et al. 2013). Since Ensmann & Burrows (1992) also assume electron/ion coupling, one would not expect a spike in temperature in their simulations either. The temperature spike in their calculations has been attributed to the simplification they use to calculate their opacities (Tolstov et al. 2013).

However, as the ions and electrons decouple, the ions no longer lose energy to the radiation field and the ion temperature can dramatically spike at the shock front. As the shock density decreases (e.g., in a SN remnant), the ions and electrons no longer couple, and our “strongly coupled” assumption breaks down. Depending on the density of the surroundings,
this decoupling can occur during the SN outburst, leading to increased X-ray emission. This may be the cause of the spike in the Klein & Chevalier (1978) calculations.

3.1. Light Curves/Spectra

We show bolometric luminosities for our Pop III CC SNe in Figure 5. Peak luminosities at shock breakout vary from $8 \times 10^{44}$ to $1.5 \times 10^{46} \text{ erg s}^{-1}$ in the z-series and from $3 \times 10^{44}$ to $5 \times 10^{45} \text{ erg s}^{-1}$ in the u-series. At early times the light curve is powered primarily by the conversion of kinetic energy into thermal energy by the shock, so for a given progenitor mass the peak luminosities increase with explosion energy. They peak at earlier times with greater $E$ because the shock reaches the surface of the star in less time. Breakout also generally happens later with more massive stars because they have greater radii in both series, with the exception of the 25 $M_{\odot}$ star because it dies with the smallest radius, as shown in Figure 1. For a given progenitor mass and $E$, u-series SNe are somewhat less luminous than z-series SNe, and this is due to a tradeoff between shock temperature at breakout and the radius of the star:

$$L = 4\pi \epsilon r^2 \sigma T^4.$$  \hspace{1cm} (11)

Here, $\sigma$ is the Stefan–Boltzmann constant, $\epsilon$ is the graybody correction to the blackbody luminosity assumed for the shock ($\epsilon \sim 0.1$), and $T$ is the temperature of the shock at the $\tau_{\text{th}} = 1$ surface, where $\kappa_{\text{th}} = 0.288$. The red z-series stars have radii 10–30 times greater than u-series stars of equal mass but their shocks have lower temperatures at breakout because they must do more $PdV$ work before crashing through the surface of the star. Their respective temperatures are evident in their spectra (11). The u40G model has $10^{-4}$ to $1.5 \times 10^{-2}$ for up to 300 days. However, because the shock cools over time, the region of its spectrum that is redshifted into the NIR in the observer frame dims below detection limits well before 1 year, as we discuss below.

As the shock expands, it cools, and its emission at later times is powered by the radioactive decay of $^{56}\text{Ni}$. The duration of this emission can be approximated by the radiation diffusion timescales in the ejecta:

$$t_d \sim \kappa^{1/2} M_{\text{ej}}^{3/4} E^{-3/4},$$  \hspace{1cm} (12)

where $M_{\text{ej}}$ is the mass of the ejecta, $\kappa$ is the average opacity of the ejecta, and $E$ is the explosion energy. The u-series SNe dim after about three months but the z-series explosions exhibit much more extended emission reminiscent of Type IIp SNe, whose progenitors are also thought to be red supergiants. In particular, the z40 series remain above $10^{32} \text{ erg s}^{-1}$ for up to 300 days. However, because the shock cools over time, the region of its spectrum that is redshifted into the NIR in the observer frame dims below detection limits well before 1 year, as we discuss below.

We show the evolution of the z40G and u40G spectra in Figure 6 (compare to Figure 2 in Tomimaga et al. 2011). Unlike with much more energetic Pop III PI SNe, the outer regions of the ejecta and the envelope are never fully ionized so bound–bound and bound–free transitions absorb most of the flux at the short wavelength limit of the spectrum. As the shock expands and cools, its surrounding envelope removes more flux at high energies in the spectrum and continuum emission at long wavelengths slowly rises. The flux at longer wavelengths increases with time because the ejecta cools and its surface area grows. From these profiles, it is clear that fitting a blackbody profile to a bolometric luminosity in order to approximate a spectrum can lead to serious overestimates of flux at short wavelengths because it erroneously reinstates luminosity that is actually removed by the envelope. This caveat is especially pertinent to detection thresholds for high-redshift SNe because flux from this region of the spectrum is a principal component of the NIR signal of the explosion in the observer frame. The collision of the shock with a realistic circumstellar envelope (which sets its temperature) and the opacity of the ejecta and
envelope also crucially shape the spectra in the source frame, and thus the NIR light curve in the observer frame.

4. POP III CC SN DETECTION THRESHOLDS

We calculate NIR light curves for our SNe with the photometry code developed by Su et al. (2011). Each spectrum is redshifted before removing the flux that is absorbed by intervening neutral hydrogen along the line of sight according to the method of Madau (1995). The spectrum is then dimmed by the required cosmological factors. We linearly interpolate the least sampled data between the input spectrum and filter curve to model the light curve in a given filter.

At every redshift for each SN, we calculate the NIR signal in JWST NIRCam filters above and below the Lyman limit to find the filter in which the SN is brightest. As in Whalen et al. (2013b), we find that the CC SNe in our study are most luminous just redward of 1216 Å in the source frame at the redshifts we consider. We show NIR light curves for z-series and u-series SNe at $z = 7$, 10, and 15 in Figures 7 and 8. The NIRCam photometry limit is AB magnitude 32, and so six of the eight z-series SNe will be visible for 20–600 days at $z = 7$ but only one of the u-series explosions, u40G, will be visible, and only for $\sim$30 days. Five of the z-series SNe will be visible at $z = 10$: 

**Figure 5.** Bolometric light curves for 17 Pop III CC SNe in the source frame. Upper left: all eight z-series light curves (red supergiant progenitors). Upper right, lower left, and lower right are u15, u25, and u40 (blue giant progenitor) light curves, respectively. (A color version of this figure is available in the online journal.)
Figure 6. Spectral evolution of the fireball. Left: z40G at $3.92 \times 10^5$ s (black), $5.55 \times 10^5$ s (blue), $1.49 \times 10^6$ s (red), $8.66 \times 10^6$ s (green), and $2.95 \times 10^7$ s (purple). Right: u40G at $1.87 \times 10^4$ s (black), $6.82 \times 10^4$ s (blue), $2.06 \times 10^5$ s (red), $1.21 \times 10^6$ s (green), and $9.69 \times 10^6$ s (purple).

Figure 7. NIR light curves for $z = 7$ Pop III CC SNe at 1.63 $\mu$m in the observer frame. Left: z-series. Right: u-series.

$z_{15B}$ for less than 1 day, $z_{15D}$ for 75 days, $z_{15G}$ for 220 days, $z_{40B}$ for 275 days, and $z_{40G}$ for over 400 days. At $z = 15$ the $z_{15G}$ SN is visible for 250 days and the $z_{40B}$ and $z_{40G}$ explosions can be seen for 300 days. The peak in emission for all the SNe advances to later times and is broader at earlier epochs because of cosmological redshifting. For a given mass, the NIR luminosity increases with the explosion energy.

Shock breakout is not visible from Earth in any of these SNe because the x-rays and hard UV in the transient are absorbed by neutral H at high redshift. The flux that is redshifted into the NIR varies much more rapidly (on timescales of $\sim 100$ days) than the bolometric luminosity (3–7 years) in the observer frame because the spectra evolve as the fireball expands. The light curves rise more rapidly than they fall, so they are most recognizable as transients in their early stages. However, given JWST survey times of 1–5 years, it will be possible to identify these events as SNe at any stage above detection threshold. Because they are much dimmer than Pop III PI SNe and because $z \sim 7$
galaxies will be more luminous than \( z \sim 10–15 \) protogalaxies, some Pop III CC SNe may be somewhat more challenging to discriminate from their host galaxies (with which they likely overlap in color–color space). However, they will still be more luminous than most galaxies of that epoch and may be much more frequent than PI SNe if baryons in some metal-free halos collapse and fragment into small swarms of stars rather than a few very massive ones, as the latest numerical simulations suggest.

The prospects for detecting Pop III CC SNe with JWST at \( z \gtrsim 10 \) are better than those for PI SNe (e.g., Hummel et al. 2012) because of their higher rates, but they will be too dim to be found in future all-sky NIR surveys such as Euclid, whose detection limit is AB magnitude 24 at 2 \( \mu \text{m} \), or WFIRST and WISH, whose photometry limits will be AB magnitude 27 at 2.2 \( \mu \text{m} \). They also will fall below the detection threshold in the \( \text{Y} \)-band centered at 1.0 \( \mu \text{m} \) for the Large Synoptic Survey Telescope (LSST) and the Panoramic Survey Telescope and Rapid Response System (Pan-STARRS), whose photometry limits will be at most AB magnitude 25 and 27, respectively. Our SPECTRUM calculations show that these SNe will be even dimmer in the optical in LSST and Pan-STARRS because of extinction by the IGM at high redshift. We note that while we have considered only Pop III CC SNe, our detection limits hold for progenitors of any metallicity. The central engine of the explosion depends mostly on the entropy profile of the inner 3–4 \( M_\odot \) of the star, which does not vary strongly with metallicity (Chieffi & Limongi 2004; Woosley & Heger 2007; see also Figure 1 in Whalen & Fryer 2012). The light curves and spectra of CC SNe at higher metallicities should therefore be bracketed by those of our compact blue giant and red supergiant Pop III progenitor stars.

5. CONCLUSION

Although direct detections of Pop III CC SNe will not probe the very first stellar populations, they will reveal the properties of stars in \( z \sim 10–15 \) galaxies and trace their formation rates and early galactic chemical evolution in general (Wise & Abel 2005; Weinmann & Lilly 2005; O’Shea et al. 2005; Tornatore et al. 2007; Whalen et al. 2008a, 2010; Trenti et al. 2009; Greif et al. 2010; Maio et al. 2011; Hummel et al. 2012; Johnson et al. 2013; Wise et al. 2012). Furthermore, because they will likely be brighter than most primitive galaxies of that era, explosions of low-mass Pop III stars (and indeed any CC explosion) will reveal the existence of primeval galaxies on the sky that might otherwise escape detection by JWST or future ground-based 30 m class telescopes. We note that while we have considered explosion energies of 0.6–2.4 \( \text{Be} \), higher energies may be possible for some CC SNe and would be visible at even higher redshifts.

There are also scenarios in which CC explosions can produce luminosities that rival those of PI explosions that could be visible at much earlier epochs. For example, if the shock collides with a dense shell ejected during a luminous blue variable (LBV) outburst prior to the SN, an extremely bright event in the UV can result that could be detected above \( z \sim 10 \) (e.g., Smith & McCray 2007; Smith et al. 2007; van Marle et al. 2010; Moriya et al. 2010, 2013; Tanaka et al. 2012). We are now studying the observational signatures of these superluminous Pop III Type IIn SNe (Whalen et al. 2013a). There may also be hypernovae, very energetic explosions of 40–60 \( M_\odot \) stars with energies of \( 10^{52} \text{erg} \) that are intermediate to those of CC and PI SNe. This kind of explosion, whose existence is inferred in part from the elemental abundances imprinted on a few of the most metal-poor stars found to date (Iwamoto et al. 2005), could be detected at redshifts that bridge those at which CC and PI SNe can be found, \( 15 < z < 20 \). We are also calculating light curves and spectra for such events.

A small number of Pop III CC events may proceed as gamma-ray bursts (GRBs; e.g., Bromm & Loeb 2006; Wang et al. 2012), driven either by the collapse of single rapidly rotating stars (Suwa & Ioka 2011; Nagakura et al. 2012) or by binary mergers with other 20–50 \( M_\odot \) stars (e.g., Fryer & Woosley 1998;
Fryer et al. 1999, 2007; Zhang & Fryer 2001). The prospects for mergers in particular have recently been strengthened by the discovery that Pop III stars sometimes form in binaries in simulations (Turk et al. 2009). Although X-rays from these events may be detected by Swift or its successors such as the Joint Astrophysics Nascent Universe Satellite (JANUS; Mészáros & Rees 2010; Roming 2008; Burrows et al. 2010), their afterglows (Whalen et al. 2008b) might also be detected in all-sky radio surveys by the Extended Very Large Array (eVLA), eMERLIN, and the Square Kilometer Array (SKA; de Souza et al. 2011). We are now studying detection limits for Pop III GRBs in a variety of circumstellar environments (Mesler et al. 2012, 2013).

Strong gravitational lensing by massive galaxies and clusters at \(z \sim 0–1\) could magnify flux from Pop III SNe, improving prospects for their detection (Ryderberg et al. 2013). The probability that flux from a Pop III SN would be boosted in an all-sky survey and its magnification depend on the event rate on the entire sky at a given redshift. We have performed preliminary Markov Chain Monte Carlo ray-tracing calculations that suggest that the probability that a \(z \sim 20\) event is lensed is \(\sim 1\%–5\%\) for flux boosts of \(2–5\). Much higher magnifications (10–300) are possible near the edges of massive clusters but the search volumes and probabilities of encountering high-\(z\) SNe are much smaller. We continue to study strong lensing of \(z \sim 20\) events, the highest redshifts ever attempted, in order to assess its potential to discover primeval SNe and galaxies.

Can later stages of the SN remnant be detected? Whalen et al. (2008c) found that most of the energy of 15 and 40 \(M_\odot\) Pop III SNe is eventually radiated away as \(H\) and \(He\) lines as the remnant sweeps up and shocks pristine gas. At lower redshifts this energy would instead be lost to fine structure cooling by metals. In either case, the emission is too diffuse, redshifted, and drawn out over time to be detected by any upcoming instruments. Also, unlike PI SNe, CC SNe deposit little of their energy into the cosmic microwave background (CMB) by inverse Compton scattering at \(z \sim 20\) (Kitayama & Yoshida 2005; Whalen et al. 2008c), and even less at lower redshifts because the density of CMB photons falls with cosmological expansion. Consequently, early populations of low-mass Pop III SNe will probably not impose photons falls with cosmological expansion. Consequently, early populations of low-mass Pop III SNe will probably not impose excess power on the CMB at small scales (Oh et al. 2003). For the same reason it will probably not be possible to directly image Sunyaev–Zeldovich fluctuations from individual Pop III CC SN remnants with the Atacama Cosmology Telescope or the South Pole Telescope.

Although their event rates make it unlikely that Pop III CC SNe will be detected in absorption at 21 cm at \(z > 10\), new calculations reveal that enough synchrotron emission from their remnants is redshifted into the 21 cm band above \(z \sim 10\) to be directly detected by the SKA (Meiksin & Whalen 2013). Somewhat more energetic hypernovae could be detected by existing facilities such as eVLA and eMERLIN. Whether by direct detection or by their footprint on cosmic backgrounds, these ancient SNe will soon open a new window on the \(z \sim 10–15\) universe.

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