Type Ia supernova diversity: white dwarf central density as a secondary parameter in three-dimensional delayed detonation models

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ABSTRACT

Delayed detonations of Chandrasekhar mass white dwarfs (WDs) have been very successful in explaining the spectra, light curves and the width–luminosity relation of spectroscopically normal Type Ia supernovae (SNe Ia). The ignition of the thermonuclear deflagration flame at the end of the convective carbon ‘simmering’ phase in the core of the WD is still not well understood, and much about the ignition kernel distribution remains unknown. Furthermore, the central density at the time of ignition depends on the still uncertain screened carbon fusion reaction rates, the accretion history and cooling time of the progenitor, and the composition. We present the results of 12 high-resolution three-dimensional delayed detonation SN Ia explosion simulations that employ a new criterion to trigger the deflagration to detonation transition (DDT). The simulations fall into three ignition categories: relatively bright SNe with five ignition kernels and a weak deflagration phase (three different central densities); relatively dim SNe with 1600 ignition kernels and a strong deflagration phase (three different central densities) and intermediate SNe with 200 ignition kernels (six different central densities). All simulations trigger our DDT criterion and the resulting delayed detonations unbind the star. We find a trend of increasing iron group element (IGE) production with increasing central density for all three categories. The total $^{56}$Ni yield, however, remains more or less constant, even though increased electron captures at high density result in a decreasing $^{56}$Ni mass fraction of the IGE material. We attribute this to an approximate balance of $^{56}$Ni producing and destroying effects. The deflagrations that were ignited at higher density initially have a faster growth rate of subgrid-scale turbulence. Hence, the effective flame speed increases faster, which triggers the DDT criterion earlier, at a time when the central density of the expanded star is higher. This leads to an overall increase of IGE production, which offsets the percental reduction of $^{56}$Ni due to neutronization.

Key words: nuclear reactions, nucleosynthesis, abundances – supernovae: general.

1 INTRODUCTION

Type Ia supernovae (SNe Ia) have come to fame as the Universe’s most luminous standardizable candles – crucial ingredients to the study of dark energy and cosmology (e.g. Riess et al. 1998; Schmidt et al. 1998). A limiting factor on the precision of using SNe Ia as distance indicators is the inherent scatter in their normalized light curves (e.g. Wood-Vasey et al. 2007). A better understanding of the intrinsic variation of supernova brightnesses and spectra is needed (e.g. Albrecht et al. 2006; Miknaitis 2007). Simulations of SN Ia explosions are already being used to aid in improving the precision of cosmological distance measurements based on supernovae in the future (e.g. Blondin, Mandel & Kirshner 2011). In addition, SNe Ia also play a critical role in galaxy gas kinematics (e.g. Scannapieco et al. 2008), positron production (e.g. Chan & Lingenfelter 1993) and chemical evolution (e.g. Matteucci & Greggio 1986). Detailed modelling of the explosions is therefore useful for understanding the origin of the Galactic 511-keV line, the origin and evolution of heavy elements, and kinetic supernova feedback and for measuring the Hubble parameter as a function of redshift. The standard model of SNe Ia relies on the nuclear fusion of the initial composition (predominantly $^{12}$C and $^{16}$O) of a massive white dwarf (WD) star to more tightly bound nuclei to power the explosion (Hoyle & Fowler 1960). The exact nature of the progenitor systems and details of the dynamics of the nuclear burning processes, however, are not known. Among the leading scenarios are the Chandrasekhar mass models, in which a WD accretes matter from a companion star and grows in mass to near the Chandrasekhar limit until pycnonuclear carbon fusion reactions (Cameron 1959) start taking place. Once

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carbon fusion reactions produce more energy than is carried away by neutrino losses, the core becomes convective, and when the nuclear burning time of a fluid element becomes shorter than the eddy turnover time a deflagration flame may be born (e.g. Woosley 1990). Numerical simulations of the convective stage leading up to the ignition of the deflagration were performed by Höflich & Stein (2002), Kuhlen, Woosley & Glatzmaier (2006), Piro & Chang (2008), Piro & Bildsten (2008) and Zingale et al. (2009). The central density of the WD decreases significantly during the simmering phase between the onset of carbon burning and the ignition of the deflagration (e.g. Lesaffre et al. 2006; Piro & Bildsten 2008). Those calculations, however, are not taking electron captures and the Urca process correctly into account, and some uncertainty in the evolution remains. The rate of the screened $^{13}$C–$^{12}$C fusion reaction is still quite uncertain (e.g. Itoh et al. 2003; Gasques et al. 2005, 2007; Jiang et al. 2007). The central density at the time of ignition, however, depends only mildly on the exact value of this reaction rate (Cooper, Steiner & Brown 2009; Iapichino & Lesaffre 2010).

More important is the initial mass and the accretion and cooling history of the WD, which determines the thermodynamic state of the interior. This results in a range of possible central densities at ignition, from less than $2 \times 10^5$ to over $5 \times 10^6 \text{g cm}^{-3}$ (Lesaffre et al. 2006).

Metallicity has a considerable impact on the supernova brightness (e.g. Timmes, Brown & Truran 2003; Travaglio, Hillebrandt & Reinecke 2005; Bravo et al. 2010). In contrast, the ignition density has been shown to depend rather weakly on metallicity and the CO ratio (Lesaffre et al. 2006). If the initial deflagration flame can transition into a detonation (e.g. Khokhlov, Oran & Wheeler 1997; Röpke 2007; Woosley 2007; Woosley et al. 2009), then good agreement of the models with observations can be obtained (e.g. Röpke & Niemeyer 2007; Bravo & García-Senz 2008; Kasen, Röpke & Woosley 2009). A successful explosion model has to reproduce the observed range of peak absolute magnitudes (i.e. $^{56}$Ni masses) and the width–luminosity relation and scatter thereof. Furthermore, the observed correlation between the brightness of an event and the delay time or age of the host stellar population has to be explained (e.g. Gallagher et al. 2008). Recently, a connection between the age of the host stellar population and the SN Ia brightness was presented in Townsley et al. (2009), Krueger et al. (2010) and Jackson et al. (2010). Under the constraints of limited computational resources, we chose six different densities for a setup with an intermediate number of ignition points (200 kernels), and three densities each for the setups with the least (5 kernels) and the most (1600 kernels) ignition points. The central densities are such that they cover the distribution of ignition densities expected from different cooling ages and accretion histories (Lesaffre et al. 2006). The ignition spark configurations are selected in a way that SNe with a range of brightnesses with $^{56}$Ni masses between $\sim$0.45 and 1.1 $M_{\odot}$ are obtained.

2 METHODS AND SIMULATIONS

The large computational demands of the high-resolution three-dimensional simulations we perform prevented a statistical framework approach similar to the one presented in Townsley et al. (2009), Krueger et al. (2010) and Jackson et al. (2010). Under the constraints of limited computational resources, we chose six different densities for a setup with an intermediate number of ignition points (200 kernels), and three densities each for the setups with the least (5 kernels) and the most (1600 kernels) ignition points. The central densities are such that they cover the distribution of ignition densities expected from different cooling ages and accretion histories (Lesaffre et al. 2006). The ignition spark configurations are selected in a way that SNe with a range of brightnesses with $^{56}$Ni masses between $\sim$0.45 and 1.1 $M_{\odot}$ are obtained.

2.1 Initial models

All simulations presented here are full star simulations performed in three dimensions. The initial stellar models are cold, isothermal ($T = 5 \times 10^5 \text{K}$) WDs in hydrostatic equilibrium with central density $\rho_c$ ranging from $1.0 \times 10^9$ to $5.5 \times 10^9 \text{g cm}^{-3}$. The composition is assumed to be 47.5 per cent $^{12}$C, 50 per cent $^{16}$O and 2.5 per cent $^{22}$Ne (to account for solar metallicity of the zero-age main-sequence progenitor) by mass homogeneously throughout the star, resulting in an electron fraction $Y_e = 0.49886$.

A strong deflagration phase leads to more energy release and hence expansion of the star. The ensuing detonation then produces less $^{56}$Ni, leading to a dimmer event. In the multispot ignition scenarios, the strongest deflagrations are obtained by placing an optimal number of ignition sparks approximately symmetrically about the
centre (García-Senz & Bravo 2005; Livne, Asida & Höflich 2005; Röpke et al. 2007b). While too few ignition sparks lead to an overall weak deflagration, too many of them lead to vigorous burning in the initial stage and thus an early expansion of the WD that suppresses burning in later stages of the deflagration (Röpke et al. 2006).

Asymmetric ignition spark distributions lead to a weaker deflagration phase and hence a brighter SN Ia (Röpke, Woosley & Hillebrandt 2007a; Kasen, Röpke & Woosley 2009). For the ignition of the deflagration, we use setups generated from a Monte Carlo based algorithm. The primary input parameters are the number of the ignition kernels and the distribution type. The details of the ignition process remain unknown. Woosley, Wunsch & Kuhlen (2004) and Wunsch & Woosley (2004) conclude from analytical models that multipoint ignition within the inner ∼150 km or so is a possible scenario. The total number and spatial distribution of the ignition spots, however, was not conclusively constrained by their models. We investigate three different sets of explosion models corresponding to different ignition scenarios. We choose configurations of 5, 200 and 1600 kernels which are spherically arranged around the centre of the WD following a Gaussian distribution in radius. The placement of kernels with a distance greater than 2.5 times a given variance σ is suppressed. For the setups with 5, 200 and 1600 kernels, we set σ = 0.6 × 10^3, 0.75 × 10^3 and 1.8 × 10^3 cm, respectively. The radius of the spherical ignition kernels is set to R_k = 10^6 cm. Finally, we impose a length-scale D_k, which the distances between the centres of the ignition kernels have to exceed. D_k is set to 1 × 10^6, 3 × 10^6 and 5 × 10^6 cm for the setups with 5, 200 and 1600 kernels, respectively. Note that for R_k > D_k, the sparks may partially overlap, which is the case for the setups with 200 and 1600 kernels. Within a given model suite (i.e. 5, 200 or 1600 kernels), the locations of the ignition sparks are only once randomly determined in the beginning; the resulting spatial realization of the ignition configuration is then kept fixed and identical in all the simulations with different central density. With these choices of ignition spark distributions we cover a large range of 56Ni masses, between ∼0.45 and 1.1 M⊙, which is consistent with normal SNe Ia (e.g. Contardo, Leibundgut & Vacca 2000; Stritzinger et al. 2006a,b).

2.2 Computational method

The reactive Euler equations are solved using a finite volume scheme based on the PROMETHEUS code by Fryxell, Müller & Arnett (1989), which is an implementation of the ‘piecewise parabolic method’ (PPM) of Colella & Woodward (1984). The grid resolution is 512 × 512 × 512 cells for all simulations. We use the expanding hybrid grid implementation of Röpke & Hillebrandt (2005a,b), with a uniform inner grid that contains the deflagration level set and a non-uniform outer grid that covers the remainder of the computational domain. Our simulation code is based on a large eddy simulation (LES) approach, which resolves the largest turbulent structures and models the turbulence on unresolved scales using a turbulence subgrid-scale model (for details see Schmidt, Niemeyer & Hillebrandt 2006a; Schmidt et al. 2006b). The code uses a comoving grid (Röpke 2005; Röpke et al. 2006) with an outer coarse grid following the WD’s expansion and an inner finer grid tracking the flame front. The flame itself is treated as a discontinuity separating fuel and ash; its propagation is tracked with the level set technique (Osher & Sethian 1988; Smiljanovski, Moser & Klein 1997; Reinecke et al. 1999). In this thin flame approximation, the energy liberated in the nuclear burning is released immediately behind the level set representing the flame surface. Since nuclear matter burned in a deflagration undergoes different burning than matter processed in a detonation, separate level set representations are used (Golombek & Niemeyer 2005; Röpke & Niemeyer 2007). Using a full nuclear reaction network in every computational cell to calculate the source terms for the hydrodynamics is currently still computationally too expensive for three-dimensional simulations. We solve this problem by tabulating the energy release as a function of fuel density. For the detonation, we use the new tables from Fink et al. (2010). A table for the energy release of the deflagration level set was calculated in a similar way.

2.3 DDT criterion

The transition from a subsonic deflagration to a supersonic detonation based on the Zel’dovich gradient mechanism (Zel’dovich et al. 1970) was introduced to SN Ia theory by Blinnikov & Khokhlov (1986) and further analysed by Khokhlov (1991a,b), Khokhlov et al. (1997) and Niemeyer & Woosley (1997). The main result of their studies was that such a transition is only possible in the turbulent deflagration stage, where large velocity fluctuations v′ lead to a mixing of cold fuel and hot ash up to a certain length-scale. These ‘hot spots’ are supposed to be the seeds of a DDT. Lisewski, Hillebrandt & Woosley (2000) pointed out that v′ must exceed 10^6 cm s⁻¹. Indeed, velocity fluctuations on this scale have already been found in three-dimensional deflagration simulations (Röpke 2007). Woosley (2007) argued that for DDTs there are specific restrictions on the burning properties deep in the distributed burning regime, which is the regime where strong turbulent flame interactions are expected. As DDTs cannot be resolved in full-star simulations, we employ a subgrid-scale model to calculate the probability of these transitions.

The details of this subgrid-scale DDT model, which is guided by the latest studies of the microscopic mechanism of DDTs in SNe Ia (Woosley et al. 2009), are described in a separate paper (Ciardi-Schoolmann & Röpke, in preparation). It accounts for the intensity of the turbulent velocity fluctuations as well as the fuel density ρ_fuel and fuel fraction X_fuel in the grid cells crossed by the flame front. If the probability P(v′ > v′_crit) to find velocity fluctuations larger than v′_crit in a specific area A_flame at the flame front exceeds a certain threshold A_crit, detonations are ignited in the grid cells which contain the largest velocity fluctuations. A_flame is defined as the part of the flame where ρ_fuel ∈ [0.6, 0.8] · 10^9 g cm⁻³ and X_fuel ∈ [0.3, 0.7]. To properly estimate this area, we take the fractal dimension of the flame front into account, which is ~2.36 (Kerstein 1988; Sreenivasan 1991; Woosley 2007). The number of ignitions is given by the ratio A_flame to A_crit. The criterion must hold at least for half of an eddy turn $t_{eddy} = L/V(L)$, where L is the turbulent integral scale and V(L) is the velocity at this scale. Following considerations of Woosley (2007), Ciardi-Schoolmann et al. (2009) and Röpke (2007), we assume $L = 10^6$ cm and $V(L) \approx (10^7\text{–}10^8)$ cm s⁻¹. We choose a constant value of $t_{eddy} = 0.005$ s in our analysis. We further follow Lisewski et al. (2000) and Röpke (2007) and define $v′_crit = 10^6$ cm s⁻¹ and $A_{crit} = 10^{12}$ cm² as our thresholds for the DDT criterion. While the details of the implementation are beyond the scope of this publication and will be presented elsewhere, we point out that this modelling approach is significantly different from simply fixing a certain DDT threshold density. Our criterion in addition requires strong local turbulent velocities. It is thus more restrictive and substantially reduces the number of DDTs.

3 RESULTS

The chosen distinct setups lead to different evolutions of the deflagration flame. In turn, the different evolutions of the flame front
The ensuing weaker expansion manifests itself in a higher central parameter for the $^{56}$Ni production (Röpke & Niemeyer 2007), and hence brightness, of a SN that explodes in the delayed detonation. Thus, it is not surprising that models with higher initial central density also have a higher central density at $t = t_{DDT}$, which is a proxy for the amount of fuel at densities high enough that it will be burned to IGE in the detonation.

We find that for all ignition kernel distributions, the total yield of IGE material within an ignition distribution suite increases with $\rho_c$ (see Fig. 5 and Table 2). The total yield of $^{56}$Ni appears flat with $\rho_c$ for the model suites with five and 200 ignition kernels; only the model suite with the strongest deflagration phase (1600 ignition kernels) has a trend of decreasing $^{56}$Ni with $\rho_c$ (see Fig. 6 and Table 2).

4 DISCUSSION

We attribute the almost constant $^{56}$Ni mass to a coincidental balance of the competing effects presented in Section 1—the larger electron capture rates at higher central density are offset by a greater total mass in IGEs due to the inherent compactness of the WD and faster evolution of the flame towards DDT. The high-density simulations with five and 200 ignition sparks exhibit more subgrid-scale energy generation at early times and therefore higher flame speeds initially. The still highly turbulent deflagration flame reaches the outer layers of the WD with low fuel density faster, and, consequently, the DDT criterion triggers earlier when the central density of the star is still lower (see Table 1). As a result, these simulations produce more

Table 1. DDT attributes for all models. Tabulated are the nuclear energy released $E_{\text{nuc}}$, the central density $\rho_c$ and the (average) density of the first DDT spot(s) $\bar{\rho}_1(t = t_{DDT})$ at the time $t_{DDT}$ when the first DDT(s) occurred.

<table>
<thead>
<tr>
<th>Model</th>
<th>$\rho_c(t = 0)$</th>
<th>$t_{DDT}$</th>
<th>$E_{\text{nuc}}(t = t_{DDT})$</th>
<th>$\rho_c(t = t_{DDT})$</th>
<th>$\bar{\rho}<em>1(t = t</em>{DDT})$</th>
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<tr>
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<tr>
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<td>0.667</td>
<td>1.706</td>
<td>0.746</td>
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<td>0.803</td>
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<td>0.718</td>
<td>0.705</td>
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<tr>
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<td>0.990</td>
<td>0.827</td>
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<tr>
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<td>0.757</td>
<td>1.087</td>
<td>0.875</td>
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</table>
but also radioactive nuclides such as $^{57}\text{Ni}$ and $^{55}\text{Co}$ contribute. Monthly Notices of the Royal Astronomical Society C mainly made up of stable iron group nuclides such as $^{58}\text{Ni}$ and $^{54}\text{Fe}$, the increased production of IGE material does not reflect in larger hind the slowly moving deflagration flame front, and consequently more neutron rich stable Fe group isotopes occur most copiously be-

ing such large amounts of strongly neutronized IGE matter cannot, however, make up most SN Ia events, due to the unusual isotopic composition (e.g. Woosley 1997).

For simulations of delayed detonations in SNe occurring via a DDT, the detonation is generally put in 'by hand' (e.g. Arnett & Livne 1994; Livne 1999; Gamezo, Khokhlov & Oran 2005; Bravo & García-Senz 2008; Krueger et al. 2010), usually by choosing a critical density where a deflagration transitions to a detonation. GARCÍA-SENZ 2008; KRUEGER ET AL. 2010), usually by choosing a critical density where a deflagration transitions to a detonation. Recently, Jackson et al. (2010) investigated the effect the particular choice of such a transition density has. They found a quadratic choice of such a transition density has. They found a quadratic dependence of the IGE yield on the log of the transition density.

Table 2. Nucleosynthetic yields for all models. Tabulated are the total WD mass $M_{\text{tot}}$ and the final masses of $^{12}\text{C}$, $^{16}\text{O}$, intermediate mass elements, IGEs and $^{56}\text{Ni}$ ($M_{\text{IC}}, M_{\text{IME}}, M_{\text{IGE}}$ and $M_{\text{SN}}$). Furthermore tabulated are the masses of IGEs and $^{56}\text{Ni}$ at the time $t_{\text{DDT}}$ when the first DDT(s) occurred ($M_{\text{IGE}}^{\text{def}}$ and $M_{\text{SN}}^{\text{def, IME}}$), as well as their respective relative fractions of the final masses, ($M_{\text{IGE}}^{\text{def}}/M_{\text{IGE}}$) and ($M_{\text{SN}}^{\text{def, IME}}/M_{\text{SN}}$).

<table>
<thead>
<tr>
<th>Model</th>
<th>$\rho_0 (t=0)$ ($10^9 \text{g cm}^{-3}$)</th>
<th>$M_{\text{tot}}$ ($M_\odot$)</th>
<th>$M_{\text{IC}}$ ($M_\odot$)</th>
<th>$M_{\text{IME}}$ ($M_\odot$)</th>
<th>$M_{\text{IGE}}$ ($M_\odot$)</th>
<th>$M_{\text{SN}}$ ($M_\odot$)</th>
<th>$M_{\text{IGE}}^{\text{def}}$ ($M_\odot$)</th>
<th>$M_{\text{SN}}^{\text{def, IME}}$ ($M_\odot$)</th>
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<td>0.449</td>
<td>0.858</td>
<td>0.463</td>
<td>0.665</td>
</tr>
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</table>

Figure 2. Shown is the mass of $^{56}\text{Ni}$ produced in the different explosions as a function of the total nuclear energy liberated during the deflagration phase up to $t = t_{\text{DDT}}$.

IGEs, and even though the $^{56}\text{Ni}$ fraction of the IGEs is lower due to increased neutronization (see Fig. 3), the total amount of $^{56}\text{Ni}$ remains roughly constant. The difference between $^{56}\text{Ni}$ and IGE is mainly made up of stable iron group nuclides such as $^{58}\text{Ni}$ and $^{54}\text{Fe}$, but also radioactive nuclides such as $^{57}\text{Ni}$ and $^{59}\text{Co}$ contribute. The model suite with 1600 ignition kernels, which has such a high density of ignition sparks that the whole central region is filled with burning products of the deflagration, behaves differently. Due to the numerous ignition sites, a large part of central mass of the star is already burned in the deflagration before the first DDT occurs. The IGEs produced in the deflagration, where most of the electron captures occur, are a large fraction of the total IGEs produced (see Table 2). The electron captures and resulting shift of $^{56}\text{Ni}$ towards more neutron rich stable Fe group isotopes occur most copiously behind the slowly moving deflagration flame front, and consequently the increased production of IGE material does not reflect in larger $^{56}\text{Ni}$ masses for cases where the deflagration contributes most of the IGE mass. The strong deflagration and vigorous expansion leads to such low central densities at $t = t_{\text{DDT}}$ that the ensuing detonation cannot produce sufficient $^{56}\text{Ni}$ to counter this trend. Events produc-

Figure 3. Shown is the relative mass fraction of $^{56}\text{Ni}$ to IGEs produced in the different explosions as a function of the total nuclear energy liberated during the deflagration phase up to $t = t_{\text{DDT}}$. It is evident that the strength of the deflagration [as measured by $E_{\text{nuc}}(t = t_{\text{DDT}})$] is a very good proxy for the mass ratio of $^{56}\text{Ni}$ to IGEs.
deflagration evolutions on the detonation initiation into account. We
point out that typical densities where our DDT criterion triggers (see
Table 1) are lower than $10^7$ g cm$^{-3}$. Jackson et al. (2010) have shown
that the variance of the $^{56}$Ni yield for a statistical set of simulations is
relatively large for such a low choice of transition density (see fig. 3
from their work), in agreement with our observed large range of $^{56}$Ni
masses obtained. Numerous other obvious differences between the
simulation sets exist, including the nature of the propagation and
the nuclear energy release of the burning fronts (level sets versus
reaction progress variables), the dimensionality of the simulations
(three dimensions versus two dimensions) or the structure of the
computational mesh (adaptive mesh refinement versus expanding
grid).

In this context, note the work of Meakin et al. (2009), who present
a suite of supernova explosion models with different offsets for the
initial deflagration bubble. Although their single-bubble off-centre
ignition scenario does not explore central density at the time of the
ignition of the deflagration as a parameter, their result that a strong
deflagration phase need not necessarily result in less $^{56}$Ni produced
is the same. They also find that the total amount of IGEs decreases
for models that had a more vigorous deflagration phase (leading
to more expansion) before the detonation is triggered, but the $^{56}$Ni
yields remains approximately constant (see fig. 12 of their work).

5 CONCLUSIONS

We have performed 12 three-dimensional hydrodynamical simula-
tions for delayed detonation SNe Ia for a range of central densities
and ignition conditions. We find a trend of increasing IGE pro-
duction with central density within each set of ignition conditions.
This is because the high central density WDs are more compact and
the flame evolves faster; the DDT occurs sooner when more un-
burned material is still above the density threshold ($\approx 10^7$ g cm$^{-3}$)
where a detonation will still produce IGEs. In spite of the larger
IGE mass, the more vigorous neutronization occurring in the high-
density models during the deflagration phase yields $^{56}$Ni masses that
are more or less constant with $\rho_c$ for the brighter SNe. Only dim
SNe, which have a strong deflagration phase and expansion prior to
the DDT, exhibit a trend of decreasing $^{56}$Ni mass with increasing
density, since the increased neutronization in the deflagration phase
cannot be compensated for by the relatively weak detonation phase.
This trend, however, is of secondary importance when compared to
the effects of varying the ignition kernel distribution. For a given
ignition kernel spatial distribution, the central density therefore in-
fluences the brightness of the supernova event only as a secondary
parameter. From the works of Townsley et al. (2009) and Bravo et al.
(2010), it appears that the same holds for composition, i.e. metal-
llicity and C/O ratio. Indeed, based on an analysis of high-quality
$V$ and $B$-band light curves of SNe Ia from the Carnegie Supernova
Project, Höflich et al. (2010) propose that the composition and cen-
tral density are two independent secondary parameters for SN Ia
light curves. In light of the importance of the ignition configuration
of the deflagration for the brightness of the SN, it is most crucial
to establish how the central density at ignition (cooling time) and
metallicity affect the statistical properties (notably number and lo-
cation) of the ignition sparks themselves, and not their respective
direct effects on the outcome of an explosion once a random ig-
nition spark distribution was chosen. One should therefore aim to
quantify which effect composition, cooling and accretion history
have on the ignition process, for example by mapping them into
the exponentiation parameter $C_e$ of the stochastic ignition prescrip-
tion of Schmidt & Niemeyer (2006). This would require a better
understanding of the physics leading up to ignition, including the nature of the convection and effects of electron captures and the convective Urca process.

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SN Ia diversity: central density variations

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