TYPE Ia SUPERNOVA EXPLOSIONS FROM HYBRID CARBON–OXYGEN–NEON WHITE DWARF PROGENITORS

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ABSTRACT

Motivated by recent results in stellar evolution that predict the existence of hybrid white dwarf (WD) stars with a C–O core inside an O–Ne shell, we simulate thermonuclear (Type Ia) supernovae from these hybrid progenitors. We use the FLASH code to perform multidimensional simulations in the deflagration-to-detonation transition (DDT) explosion paradigm. Our hybrid progenitor models were produced with the MESA stellar evolution code and include the effects of the Urca process, and we map the progenitor model to the FLASH grid. We performed a suite of DDT simulations over a range of ignition conditions consistent with the progenitor’s thermal and convective structure assuming multiple ignition points. To compare the results from these hybrid WD stars to previous results from C–O WDs, we construct a set of C–O WD models with similar properties and similarly simulate a suite of explosions. We find that despite significant variability within each suite, trends distinguishing the explosions are apparent in their 56Ni yields and the kinetic properties of the ejecta. We compare our results with other recent work that studies explosions from these hybrid progenitors.

Key words: hydrodynamics – nuclear reactions, nucleosynthesis, abundances – supernovae: general – white dwarfs

1. INTRODUCTION

Type Ia supernovae (SNe Ia) are bright stellar explosions that produce ~0.6 $M_\odot$ of radioactive 56Ni, the decay of which powers the light curve and leads to a relation between the peak brightness of an event and the rate of its decline from maximum (Phillips 1993). This relation enables SNe Ia to be used as “standard candles” for cosmological studies, and this use led to the discovery that the expansion of the universe is accelerating due to dark energy (Riess et al. 1998; Perlmutter et al. 1999; Leibundgut 2001).

Despite intense study (driven in part by their application as distance indicators for cosmology), we still have only an incomplete understanding of the explosion mechanism, and fundamental questions, such as the likely progenitor system(s), persist. It is widely accepted that SNe Ia result from the thermonuclear explosion of a white dwarf (WD) composed largely of C and O, with this understanding going back many decades (Hoyle & Fowler 1960; Arnett et al. 1971). The rapid thermonuclear fusion of C and O in a WD releases enough energy to unbind it, produces the 56Ni necessary to power the light curve, and explains the lack of H observed in the spectra.

There are, however, several possible progenitor systems for such a configuration. All models involve a binary system and at least one C–O WD, which follows from the original association of SNe Ia with C–O burning under degenerate conditions (Hoyle & Fowler 1960). Some proposed systems posit a single WD, the single-degenerate (SD) paradigm, and some posit the merger or collision of two WDs, the double-degenerate (DD) paradigm, and within these are variations.

The “classic” model is the Chandrasekhar-mass model, in which a WD gains mass from a companion, a main-sequence or red giant star, or perhaps an He WD, and a thermonuclear runaway occurs just as it approaches the Chandrasekhar limiting mass (Hoyle & Fowler 1960; Truran & Cameron 1971; Whelan & Iben 1973; Nomoto et al. 1984). Alternately, in the sub-Chandrasekhar-mass scenario, explosive burning in the accreted layer triggers a detonation at the surface or in the core of a lower-mass WD (Nomoto 1980; Woosley et al. 1980; Sim et al. 2010).

The other main class of models is the DD progenitor (Webbink 1984; Iben & Tutukov 1984), in which two WDs inspiral and merge. This scenario includes inspiraling pairs, collisions, violent mergers, and also the “core-degenerate” model, where the merger takes place in a common envelope (Raskin et al. 2009; Kashi & Soker 2011; Pakmor et al. 2011, 2012; Shen et al. 2012). Also see Hillebrandt & Niemeyer (2000), Howell (2011), Hillebrandt et al. (2013), and Calder et al. (2013) for additional discussion.

The observational evidence of one progenitor versus another is conflicting. SNe Ia show a wide range of luminosities and also the possibility that there are two classes of progenitor (Mannucci et al. 2006; Howell et al. 2009; Howell 2011). Observational and population synthesis (Yungelson & Livio 2000; Maoz 2008; Toonen et al. 2012; Claeys et al. 2014) arguments suggest that there simply may not be enough Chandrasekhar-mass progenitors to explain the observed SN Ia rate. There is, however, disagreement over the significance of these observations (Hachisu et al. 2008), and our understanding of the different progenitor systems is subject to uncertainties in aspects of their pre-SN evolution, including the effects of self-absorption on the number of supersoft X-ray sources predicted from the SD channel (Nielsen et al. 2013). Certainly there is disagreement in the interpretation of observations that stems from uncertainty in the accretion phase of SD evolution (Hachisu et al. 2010).

Additionally, the oft-cited claim that the WD in the SD channel would lose mass via nova explosions, thereby preventing it from reaching the Chandrasekhar mass, is questioned on the basis of high observed WD masses in...
cataclysmic variables (CVs) relative to pre-CVs (Zorotovic et al. 2011). Prior work on novae and rapidly accreting WDs strongly suggests that novae will not be able to grow (Denissenkov et al. 2013a, 2014), especially not from WD masses of 0.83 $M_\odot$, suggested to be the mean mass of WDs in CVs in Zorotovic et al. (2011). Recently, binary population synthesis calculations by Wijnen et al. (2015) explored both thermal timescale mass transfer and mass increase during novae as potential solutions, finding that both mechanisms result in contradictions with observations and cannot satisfy the observed WD mass difference.

However, if it is possible to get an SN Ia out of the SD scenario, then “hybrid” C–O–Ne WDs (WDs with a C–O core in an O–Ne shell; Denissenkov et al. 2015) may play a key role (see Section 2). Prior stellar evolution work by García-Berro et al. (1997) produced an O–Ne WD with a $^{12}$C abundance as high as 0.048 and suggested that this is sufficient to ignite an explosion if the WD accreted matter to approach the Chandrasekhar mass. The hybrid WD of Denissenkov et al. (2015) has a significantly higher average $^{12}$C abundance of 0.17, and such WDs would provide readily ignitable SN Ia progenitors that are already very close to the Chandrasekhar limiting mass and are therefore perhaps the most likely to produce a SN Ia.

Acknowledging the ignitability of hybrid C–O–Ne WDs, Wang et al. (2014) and Meng & Podsiadlowski (2015) presented binary population synthesis studies for these hybrid WDs and suggested that they find consistency with the observed properties of Type Iax events, though uncertainties remain in estimating the rate of SNe Ia from hybrid WDs due to the uncertain carbon-burning rate and common-envelope ejection efficiency. The study we present here addresses SN Ia explosions from these hybrid models, and other research groups have also recently explored this possibility using hydrodynamic simulations. Kromer et al. (2015) performed pure deflagration simulations from C-core models, and Bravo et al. (2016) simulated explosions from a variety of progenitor models, including WDs with similar stellar evolution to those of Denissenkov et al. (2015). We compare these explosion studies to ours in the discussion below.

While there is uncertainty, some contemporary observations do strongly support the SD progenitor. Events like PTF11kx and others show distinct circumstellar shells of material that can be best explained in the SD context (Dilday et al. 2012; Silverman et al. 2013). The supernova remnant 3C 397 is heralded as a case where only an explosion from a Chandrasekhar-mass progenitor can produce the nuclei seen in the remnant, due to the need for electron captures at high density (Yamaguchi et al. 2015). The recent observation of a UV pulse (Cao et al. 2015) in the early evolution of an SN Ia also supports the SD model. Observations of remnants also offer support for Chandrasekhar-mass explosions, including wind-blown shells in RCW 86 (Williams et al. 2011) and shocked circumstellar material/bubble in the Kepler remnant (Chiotellis et al. 2012; Burkey et al. 2013). Altogether, there is substantial evidence that suggests that the SD channel plays a role in at least some of the observed SNe Ia (Baron 2014).

The sub-Chandrasekhar-mass model does not have the population synthesis arguments working against it, and we know that low-mass WDs in binary systems exist. Systems for which it is believed that they will evolve to an explosion have been observed (Kilic et al. 2014), potential events have been identified (Geier et al. 2013; Inserra et al. 2015), and the Type Iax subclass of SNe Ia (Foley et al. 2013; Wang et al. 2013) has been suggested as consisting of sub-Chandrasekhar-mass events themselves.

Observational evidence also supports the DD progenitor system, and the scenario is increasingly seen as the likely progenitor of some events. SN 2011fe has been intensely observed and does not show features in its spectra that would be expected if there were a normal stellar companion (Graham et al. 2015), suggesting a DD system. Super-Chandrasekhar-mass explosions like SN 2007if (Scalzo et al. 2010; Yuan et al. 2010) and SNLS 03D3bb (Howell et al. 2006) also suggest mergers. There are also many population synthesis arguments in favor of mergers as well (see Maoz et al. 2014, for a review).

1.1. The Chandrasekhar-mass SD Scenario

In the Chandrasekhar-mass scenario, the central temperature and density of the WD increase as it accretes mass from a binary companion and approaches the limiting Chandrasekhar mass. As the mass approaches the limit, central conditions become hot enough for carbon fusion to begin (via the $^{12}$C-$^{12}$C reaction), driving the development of convection throughout the interior of the WD (Baraffe et al. 2004; Woosley et al. 2004; Wunsch & Woosley 2004; Kuhlen et al. 2006; Nonaka et al. 2012). As the central temperature reaches $\sim 7 \times 10^8$ K, the fuel in a convective plume burns to completion before it can cool via expansion (Nomoto et al. 1984; Woosley et al. 2004), and a flame is born.

The nature of this burning, be it a supersonic detonation or subsonic deflagration, largely determines the outcome of the explosion. It has been known for some time that a purely supersonic burning front cannot explain observations because the supersonic front very rapidly incinerates the star without it having time to react and expand (Arnett et al. 1971). The lack of expansion allows most of the star to burn at high densities, which produces an excessive $^{56}$Ni yield and does not match the stratified composition of observed remnants (Mazzali et al. 2008). Instead, a subsonic deflagration must ignite, which allows the outer layers of the star to expand ahead of the burning front. In this case, the density of the expanding material decreases, which leads to incomplete burning of more mass and thus increased production of intermediate-mass elements. This deflagration must accelerate via instabilities and turbulent interaction, a topic that has been explored extensively in the past (Khokhlov 1993, 1995; Bychkov & Liberman 1995; Niemeyer & Hillebrandt 1995; Khokhlov et al. 1997; Cho et al. 2003, pp. 56–98; Röpke et al. 2003, 2004; Bell et al. 2004; Zingale et al. 2005; Schmidt et al. 2006a, 2006b; Zingale & Dursi 2007; Aspden et al. 2008; Ciaraldi-Schoellmann et al. 2009, 2013; Woosley et al. 2009; Hicks & Rosner 2013; Jackson et al. 2014; Hicks 2015; Poludnenko 2015).

A deflagration alone will not produce an event of normal brightness and expansion velocity (Röpke et al. 2007). Instead, the initial deflagration must transition to a detonation after the star has expanded some in order to produce abundances and a stratified ejecta in keeping with observations (Khokhlov 1991; Hoflich et al. 1995). The physics of this “deflagration-to-detonation transition” (DDT) are not completely understood, but there has been considerable study based on mechanisms involving flame fronts in highly turbulent conditions.
1.2. Systematic Effects

Contemporary observational campaigns typically investigate how the brightness and rates of SNe correlate to properties of the host galaxy such as mass and star formation rate (see Graur & Maoz 2013; Graur et al. 2015). Of particular interest is the delay-time distribution (DTD), the SN rate as a function of time elapsed from early, rapid star formation in the host galaxy, and how it may be used to constrain progenitor models (Hachisu et al. 2008; Bianco et al. 2011; Conley et al. 2011; Graur et al. 2011; Howell 2011; Maoz et al. 2012). See also the review by Maoz & Mannucci (2012). Very recent results indicate evolution of the UV spectrum with redshift, providing evidence for systematic effects with cosmological time (Milne et al. 2015).

Motivated by this interest in correlations between properties of the host galaxy and the brightness and rate of events, earlier incarnations of our group performed suites of simulations in the DDT scenario with a modified version the FLASH code (described below) to explore systematic effects on the brightness of an event measured by the yield of $^{56}$Ni (Jackson et al. 2010; Krueger et al. 2010, 2012). The study we present here explores how explosions following from a new class of “hybrid” progenitors (Denissenkov et al. 2013b, 2015; Chen et al. 2014) compare with these previous results.

2. HYBRID PROGENITOR MODELS

Rumors that the structure and evolution of stars are a solved problem (Hansen et al. 2004) are greatly exaggerated. Recent developments obtained with the modern software instrument MESA (Paxton et al. 2011, 2013, 2015) indicate that convective boundary mixing (CBM) in the cores of super asymptotic giant branch (super-AGB) stars plays a more critical role than previously thought. There are several examples in which the inclusion of CBM improves agreement between models and observations, including that of Denissenkov et al. (2013b), which studied WD interior shell convection, and Herwig (2005) and Werner & Herwig (2006), which treated He shell burning in AGB stars. Denissenkov et al. (2013b) and Chen et al. (2014) found that in some super-AGB stars, CBM halts the progression of carbon burning into the stellar core, leaving an unburned C-rich core as large as 0.2 $M_\odot$ surrounded by an O–Ne-rich intershell region extending out to the accretion layer at the end of hydrostatic carbon burning. This effect of C-flame quenching via CBM is also confirmed by the extensive parameter study on C burning in super-AGB stars in

(Blinnikov & Khokhlov 1986; Woosley 1990; Khokhlov 1991; Hofflich et al. 1995; Hofflich & Khokhlov 1996; Khokhlov et al. 1997; Niemeyer & Woosley 1997; Hofflich et al. 1998; Niemeyer 1999; Gamezo et al. 2005; Röpke 2007; Poludnenko et al. 2011; Ciaraldi-Schoolmann et al. 2013; Poludnenko 2015). These models generally reproduce the observations under certain assumptions about the ignition (Townsley et al. 2009), but research has shown that the results are very sensitive to the details of the ignition (Plewa et al. 2004; Gamezo et al. 2005; García-Senz & Bravo 2005; Röpke et al. 2007; Jordan et al. 2008). In our simulations, we initialize a detonation once the deflagration front reaches a characteristic DDT fuel density, which controls the degree of expansion the star undergoes during the deflagration stage. The implementation details are described further in Section 3.4.

Figure 1. Abundance profile of MESA progenitor (MESA) and its reconstruction on a uniform grid at 4 km spatial resolution with the hydrostatic equilibrium condition of Equation (2) enforced (FLASH). The reduced set of nuclides are shown, where $X_{\text{Ne}} = 1 - X_{\text{C}} - X_{\text{Ne}} - X_{\text{C}}$ for the abundances labeled FLASH. Solid lines and circles denote the abundances used in FLASH, whereas plain solid lines denote the abundances used in MESA.

Farmer et al. (2015). Nevertheless, the true test of whether or not CBM is effective enough to quench the carbon flame awaits sufficiently resolved hydrodynamic simulations or further evidence via observational consequences.

Accounting for CBM, Denissenkov et al. (2015) explored the stellar evolution of a super-AGB star with initial mass of 6.9 $M_\odot$ and obtained the hybrid WD that is the focus of the present work. After hydrostatic carbon burning has ceased, the WD accretes carbon-rich material at its surface, leading to the rise of temperature near its center. This results in carbon burning in the upper layer of the small carbon-rich core, which, together with the thermal effects of the $^{22}$Ne/$^{23}$Na Urca process, provides off-center heating (Denissenkov et al. 2015) that drives convection throughout the entire WD except the carbon-rich core.

Convection subsequently mixes the carbon-poor material in the O–Ne intershell region with carbon-rich material on the accreted layer and also partially mixes carbon-rich material from the core with the carbon-poor material in the O–Ne intershell. This proceeds along with accretion and carbon burning, until the latter yields peak temperatures near 10$^6$ K, around which the local heating time is shorter than the eddy turnover time so as to ignite thermonuclear runaway (Wunsch & Woosley 2004). At this point, the carbon-rich core has been significantly depleted of carbon and consists mostly of $^{16}$O and $^{20}$Ne, while the O–Ne intershell region has been enriched to a $^{12}$C abundance of $\approx 0.14$ due to convective mixing. This scenario, immediately preceding the SN Ia-like explosion, is shown in Figures 1 and 2. The process of mapping this MESA progenitor into hydrostatic equilibrium (HSE) in FLASH is described in Section 3.2.

This hybrid WD has the interesting property that its mass before the onset of accretion is 1.06 $M_\odot$, naturally closer to the Chandrasekhar limit than a traditional C–O WD. Such hybrid WD progenitors thus require less mass accretion to approach the Chandrasekhar limit, which helps to resolve one of the difficulties with the SD progenitor system (Chen et al. 2014; Denissenkov et al. 2015; Kromer et al. 2015). The mass of this hybrid WD following accretion is 1.36 $M_\odot$. 
We note that these models include the influence of the Urca process on the convective phase of the pre-explosion progenitor. Our progenitor profiles are taken directly from MESA models presented in Denissenkov et al. (2015). These include contributions to the energy from thermal energy produced by the Urca process, but the underlying mixing length theory was not modified correspondingly. Thus, the effect on the convection is only due to the energy loss/generation rate. Our progenitor profiles are shown in Figures 1 and 2 and correspond to the models in Figure 9(a) of Denissenkov et al. (2015). In this regard, this progenitor differs from the carbon-core models of Kromer et al. (2015) in that it includes the pre-explosion convective burning phase that spreads the carbon enrichment throughout the star before ignition of the flame front. While the progenitor we study consists mostly of $^{16}$O and $^{20}$Ne, having the average composition of $^{12}$C = 0.17, $^{16}$O = 0.42, $^{20}$Ne = 0.32, it differs from the O–Ne WD of Marquardt et al. (2015) by having a much higher abundance of $^{12}$C due to accumulation and mixing of $^{12}$C material during the accretion phase, as described above. As discussed below in Sections 3.1 and 4.2, given the temperature profile of Figure 2, this available $^{12}$C is sufficient to drive both a subsonic deflagration and subsequently a supersonic detonation front as in previous work that applied the same DDT approach to C–O WD progenitors (Krueger et al. 2012).

3. METHODOLOGY

A few significant new developments in our computational methods were necessary to simulate the explosion of the hybrid C–O–Ne WD. In Section 3.1, we obtain the steady–state detonation structure for the hybrid C–O–Ne fuel and compare its detonation characteristics with those of C–O fuel to analyze the suitability of our combustion model in FLASH. Then in Section 3.2 we map the hybrid C–O–Ne WD into a uniform spatial grid to initialize FLASH while taking care to preserve HSE. We comment in Section 3.3 on our FLASH combustion model and in Section 3.4 on the DDT scheme. Finally, we describe the simulation geometry and the adaptive mesh refinement used in Section 3.5.

3.1. Modifications for C–O–Ne Burning

The combustion model in FLASH that we use for SN Ia simulations (Calder et al. 2007; Townsley et al. 2007, 2009, 2016; Seitenzahl et al. 2009; Jackson et al. 2014) separates the burning into four states: unburned fuel, C-fusion ash, a silicon-group-dominated nuclear statistical quasi-equilibrium (NSQE) state, and a full nuclear statistical equilibrium (NSE) state dominated by iron-group elements (IGEs). The progress of combustion from one of these states to the next is tracked by three scalar progress variables whose dynamics is calibrated to reproduce the timescales of reactions that convert material among these states. Previous work has focused on fuel mixtures composed principally of $^{12}$C, $^{16}$O, and $^{22}$Ne. Simulation of the hybrid models required extension of this burning model to account for the presence of $^{20}$Ne as a large abundance in the fuel. Here we describe both how $^{20}$Ne is processed during combustion and the modifications made to the burning model in order to incorporate $^{20}$Ne burning.

The burning stages above are determined by the hierarchy of timescales for the consecutive consumption of C and O via fusion and Si via photodisintegration and alpha capture. Investigation of the inclusion of $^{20}$Ne focused on whether an additional Ne-consumption stage would be required, and, if not, what stage should include Ne consumption. In order to characterize the physical burning sequence that we want to model, we performed a series of simulations of detonations propagating through WD material with the TORCH nuclear reaction network software (Timmes 1999, 2015).\(^6\) TORCH is a general reaction network package capable of solving networks with up to thousands of nuclides. A mode is implemented that computes the one-dimensional spatial thermodynamic and composition structure of a steady-state planar detonation using the Zel’dovich, von Neumann, and Döring (ZND) model (Fickett & Davis 1979; Townsley et al. 2016). We use a reaction network composed of 225 nuclides consisting of the 200 nuclides in Woosley & Weaver (1995) in addition to the 25 neutron-rich nuclides added by Calder et al. (2007) to improve coverage of electron capture processes in the Fe group.

For the multi-species fuel and ash relevant to typical WD material, the ZND detonation exhibits the stages that motivate the combustion model. Figure 3 shows these stages as they appear in a ZND detonation calculation in fuel with the fractional composition ($^{12}$C = 0.50, $^{16}$O = 0.48, $^{22}$Ne = 0.02) corresponding to the composition found in the interior of a C–O WD, and a fuel density of $10^7$ g cm$^{-3}$. The evolution of the mass fractions in time following the passage of the shock through the zone of material is plotted below the density structure in Figure 3. The time ranges of the four states representing the burning stages are indicated by colors in the top panel of Figure 3, with the consecutive states separated by the $^{12}$C–$^{28}$Si, $^{16}$O–$^{28}$Si, and $^{28}$Si–$^{54}$Fe crossing times. It can be seen that the times of the density plateaus are directly comparable to the times at which the primary energy release transitions from one fuel source to another, as, e.g., when the $^{12}$C fraction has fallen to $\approx$1% of its initial value just before 10$^{-8}$ s.

In our combustion models for C–O progenitors, consumption of the initial fuel is modeled as a two-step process. The two

\(^6\) TORCH is available from http://cococubed.asu.edu, and the modified version used for this study is available from http://astronomy.ua.edu/townsley/code.
stages represent the consumption first of $^{12}\text{C}$, then of $^{16}\text{O}$, mimicking the sequence seen in the detonation structure shown in Figure 3 for “Fuel” and “Ash” stages. At the end of this second stage the material is in a $^{28}\text{Si}$-dominated NSQE state (Calder et al. 2007; Townsley et al. 2007). To determine how the burning stages change with the inclusion of $^{20}\text{Ne}$, as is the case in the hybrid C–O–Ne progenitor, we perform ZND calculations with admixture of $^{20}\text{Ne}$ ranging from 0.01 to 0.45, at the expense of $^{12}\text{C}$ content. For each composition we find the minimally overdriven solution, as was done in the C–O case. Out to the minimum in density, this solution is the same as the eigenvalue ZND solution, which corresponds to a self-supported detonation (Fickett & Davis 1979; Townsley et al. 2016). This computation gives the resolved detonation structure in C–O–Ne hybrid WD matter.

The eigenvalue detonation speeds from ZND calculations in material with $^{12}\text{C}$ fraction varying from 0.05 to 0.5 are shown in Figure 4, demonstrating that self-supported detonations in this progenitor are feasible with only small variation in speed across this range of $^{12}\text{C}$ fractions. Figure 5 shows the effect of simultaneously adding $^{20}\text{Ne}$ and reducing $^{12}\text{C}$ on the density profile. Lowering the $^{12}\text{C}$ fraction weakens the shock and lengthens the timescales of the step features, corresponding to more slowly burning fuel and ash, as might be expected from the lower energy release afforded by the $^{20}\text{Ne}$. However, it is noteworthy that no qualitatively new features arise from the change in fuel source that would suggest that more than four representative burning stages are needed.

Qualitative similarity between C–O and C–O–Ne detonation structures is visible in the mass fractions as well. To demonstrate this, the abundance structures with initial $^{12}\text{C}$ abundances of 0.5, 0.3, 0.15, and 0.05 are shown in Figure 6. We find that the $^{20}\text{Ne}$ burns simultaneously with whatever $^{12}\text{C}$ is present, producing primarily $^{28}\text{Si}$. The Ne–C-burning stage is then followed by $^{16}\text{O}$ burning to silicon-group NSQE elements and then on to NSE, just as in a model with no initial $^{20}\text{Ne}$, except for the progressively later $^{16}\text{O}$ burning time.

A graphic representation of the most significant nuclides by mass fraction and the stage in which they are important is shown in Figure 7 for an initial $^{12}\text{C}$ fraction of 0.15, representative of the majority of the hybrid stellar profile (see Figure 1). Nuclides are categorized based on the time at which they were maximally abundant in the network and shaded by their maximum abundance. Nuclides within the purple color palette were maximally abundant during the initial “fuel-burning” stage after the beginning of fusion and before the $^{28}\text{Si}$ becomes equally abundant with $^{12}\text{C}$, the $^{12}\text{C}$–$^{28}\text{Si}$ crossing time $t_{\text{fa}}$. Blue nuclides were maximally abundant after $t_{\text{fa}}$ but prior to the $^{16}\text{O}$–$^{28}\text{Si}$ crossing time $t_{\text{aq}}$, and nuclides maximally abundant after $t_{\text{aq}}$ but before the $^{28}\text{Si}$–$^{54}\text{Fe}$ crossing time $t_{\text{qf}}$ are plotted in green. Nuclides maximally abundant following
are IGEs together with protons and alpha particles and are shaded in orange. Note that for the timescales we adopt the notation of Townsley et al. (2016), which describes the burning stages in detail.

The burning model we are using computes the rates for progression through the burning stages from the local temperature and the energy release from a set of major fuel abundances, previously including $^{12}$C, $^{16}$O, and $^{22}$Ne. We have found here that any $^{20}$Ne is consumed along with $^{12}$C and that otherwise the burning is quite similar to that with just $^{12}$C and $^{16}$O as principal constituents. In consequence, the only necessary modification to our burning model is to include $^{20}$Ne in the abundances of the initial state in the burning model. This accounts for the difference in binding energy of $^{20}$Ne compared to $^{12}$C and gives lower burning temperatures. Additionally, throughout the majority of the progenitor outside the core ignition region, due to prior convective mixing, the $^{12}$C content is high enough that we will extrapolate the laminar flame speeds of Timmes & Woosley (1992) and Chamulak et al. (2008). We consider this a reasonable approximation since much of the flame propagation is dominated by Rayleigh–Taylor overturn and turbulence.

3.2. Mapping a MESA Profile to FLASH While Preserving HSE

The temperature and composition at the base of the convective zone in the hybrid C–O–Ne WD provide a natural flame initialization region for simulation of thermonuclear runaway in the DDT scenario, so we map the MESA profile into the FLASH domain, preserving its features at 4 km spatial resolution. We do this by first converting the MESA model to a uniform grid by mass-weighted averaging of quantities in MESA zones with spacing less than 4 km and using quadratic interpolation to estimate quantities where MESA zones have spacing greater than 4 km. Although our combustion model in FLASH does not evolve nuclide abundances, it uses the initial abundances of $^{12}$C, $^{20}$Ne, and $^{22}$Ne (assuming that the rest is $^{16}$O) to compute the initial mean nuclear binding energy and electron fraction. Therefore, we also represent the full set of

Figure 6. Mass fraction evolution for ZND detonations with varying initial $^{12}$C and $^{20}$Ne mass fractions, calculated for an initial density of $10^7$ g cm$^{-3}$.

Figure 7. Our reaction network of 225 nuclides is shown, with nuclides shaded by their maximum abundance and categorized based on whether they were maximally abundant before the earliest of the following crossing times: $^{12}$C–$^{28}$Si (purple), $^{16}$O–$^{28}$Si (blue), $^{28}$Si–$^{56}$Fe (green). Nuclides maximally abundant after the $^{28}$Si–$^{56}$Fe crossing time are shown in orange. The initial composition used is $X_{^{12}C} = 0.48$, $X_{^{20}Ne} = 0.02$, $X_{^{16}O} = 0.15$, and $X_{^{22}Ne} = 0.35$. $^{12}$C, $^{16}$O, $^{20}$Ne, $^{22}$Ne, $^{28}$Si, and $^{28}$Si are IGEs together with protons and alpha particles and are shaded in orange. Note that for the timescales we adopt the notation of Townsley et al. (2016), which describes the burning stages in detail.
nuclides in the MESA profile by this reduced set of four nuclides in the uniformly gridded profile, requiring the carbon mass fractions to be identical because there is still sufficient $^{12}\text{C}$ in the star to sustain a detonation front. In addition, we use $^{22}\text{Ne}$ in the reduced set to account for the $Y_e$ of the full set of nuclides, and we constrain $^{20}\text{Ne}$ and $^{16}\text{O}$ to be in the same ratio $R$ in both sets of abundances. These constraints provide the following definitions for the reduced abundances used for FLASH:

\begin{equation}
X_{\text{FLASH}}^{\text{IC}} = X_{\text{MESA}}^{\text{IC}},
\end{equation}

\begin{equation}
X_{\text{FLASH}}^{\text{Ne}} = 22 \times \left( \frac{1}{2} - Y_{\text{MESA}}^{\text{Ne}} \right),
\end{equation}

\begin{equation}
R = X_{\text{MESA}}^{\text{Ne}} / X_{\text{MESA}}^{\text{O}},
\end{equation}

\begin{equation}
X_{\text{FLASH}}^{\text{IC}} = \frac{1 - X_{\text{FLASH}}^{\text{IC}} - X_{\text{FLASH}}^{\text{Ne}}}{R + 1},
\end{equation}

\begin{equation}
X_{\text{FLASH}}^{\text{Ne}} = R \times X_{\text{FLASH}}^{\text{IC}}.
\end{equation}

Nothing constrains the resulting uniformly gridded profile to be in HSE, however, so we next construct an equilibrium profile by applying the HSE pressure constraint

\begin{equation}
P_{\text{EOS}}(\rho_i, T_i, X_i) = \frac{g(\rho_i + \rho_{i-1})}{2}.
\end{equation}

Starting at the uniformly gridded profile’s central density and using its temperature and composition in each zone together with the Helmholtz equation of state (EOS) of Almgren et al. (2010), we solve Equation (2) for the density in each zone. In Equation (2), $\Delta \rho$ indicates the zone width, $T_i$ is the temperature in zone $i$, $X_i$ is the composition in zone $i$, and $g$ is the gravitational acceleration at the boundary of zone $i$ and $i-1$ due to the mass enclosed by zone $i$ and below. The resulting uniformly gridded equilibrium profile (FLASH) and the original profile (MESA) are shown for comparison in Figures 1 and 2. The values of the total mass before and after this procedure differ by $6 \times 10^{-3} M_\odot$. This procedure produced a structure that was stable in FLASH, with fluctuations in central density less than 3%, for at least 5 s with no energy deposition. Finally, we replaced the EOS routine in the public FLASH distribution with that from CASTRO (Timmes & Swesty 2000; Almgren et al. 2010) to obtain a more consistent tabulation.\footnote{The Helmholtz EOS table used in CASTRO is in the public BoxLib Microphysics repository at https://github.com/BoxLib-Codes/Microphysics.git (commit hash 45ed859b61cde8d80b31d97289868d6afde1ddc).}

3.3. Combustion and Explosion Mechanisms in the FLASH Code

To simulate the explosion from either hybrid or traditional progenitor models, we used a modified version of FLASH,\footnote{FLASH is available from http://flash.uchicago.edu, Our modifications for C-O WDs are available from http://astronomy.ua.edu/townsley/code.} an Eulerian adaptive mesh compressible hydrodynamics code developed by the ASC/Alliances Center for Astrophysical Thermonuclear Flashes at the University of Chicago (Fryxell et al. 2000; Calder et al. 2002). While FLASH is capable of evolving thermonuclear reaction networks coupled to the hydrodynamics, in order to treat a $\lesssim 1 \text{cm}$ flame front in full-star simulations of SNe Ia, we use a coarsened flame model that uses an advection-diffusion-reaction scheme to evolve a scalar variable representing the progression of burning from fuel to ash compositions as detailed in Calder et al. (2007) and Townsley et al. (2007, 2009). An additional scalar represents the burning progress from ash to intermediate-mass silicon-group elements in NSQE. A final scalar represents the burning of silicon-group elements to IGEs in NSE. The timescales for evolving these scalars are density and temperature dependent and are determined from self-heating and steady-state detonation calculations with a 200+–species nuclear reaction network (Calder et al. 2007; Townsley et al. 2016). For our evaluation of the suitability of this burning scheme for fuel with an admixture of $^{20}\text{Ne}$, see Section 3.1.

The two-dimensional (2D) models we present do not utilize a subgrid-scale model for the turbulence–flame interaction (TFI). Earlier work produced subgrid-scale TFI models by extending terrestrial combustion models to consider astrophysical flames (Schmidt et al. 2006a, 2006b; Jackson et al. 2014), however, the resulting subgrid-scale models only capture TFI in three-dimensional (3D) simulations. 2D models inherently lack true turbulence, so instead we use the minimal enhancement based on the Rayleigh–Taylor strength introduced by Townsley et al. (2007). This assumes that the TFI will self-regulate on resolved scales so that results are insensitive to the detailed treatment of the TFI. We believe this to be sufficient for 2D comparative studies like that undertaken here.

3.4. Deflagration-to-Detonation Transition

This study presents simulations of the thermonuclear explosions of both a hybrid C–O–Ne WD progenitor (Denisenkov et al. 2015) and a C–O WD progenitor similar to that used in previously published suites of SN Ia simulations (Krueger et al. 2012). In both cases, we initialize the simulations with a “matchhead” consisting of a region near or at the WD’s center that is fully burned to NSE. The energy release from this initial burn ignites a subsonic thermonuclear flame front that buoyantly rises and partially consumes the star while the star expands in response. As detailed in Townsley et al. (2009) and Jackson et al. (2010), in order to effect a DDT, we suppose the DDT point to be parameterizable by a fuel density $\rho_{\text{DDT}}$ at which the subsonic flame reaches the distributed burning regime where the flame region has become sufficiently turbulent that a supersonic detonation front may arise, self-supported by the energy release from the nuclear burning proceeding behind the detonation shock front.

We thus use a similar DDT parameterization for our C–O and hybrid C–O–Ne simulations as in the C–O WD SN Ia simulations of Krueger et al. (2010, 2012) for consistency. Our choice is to ignite the detonation when a rising plume reaches the threshold density, but reality may be quite different. For example, the turbulent intensity may be greatest at the underside of a rising plume, which would imply igniting a detonation deeper in the core and possibly later, both of which would affect the yield. For purposes of this paper, which was to compare the yield from a hybrid model to a traditional C–O model, we prefer to use a simply consistent condition across cases. Jackson et al. (2010) address the dependence of the DDT density on composition separately.

Consequently, when the deflagration reaches a point where it is at $\rho_{\text{DDT}} = 10^{7.2} \text{ g cm}^{-3}$, we place a region with its fuel stage (representing $^{12}\text{C}$ and $^{20}\text{Ne}$, if present) fully burned 32 km...
radially outward from this point that is of size 12 km in radius. Multiple DDT points may arise, but they are constrained to be at least 200 km apart. Our choice of DDT density is slightly high in order to ensure robust ignition of a detonation shock in all cases. We note that this choice strongly influences the $^{56}$Ni yield because it determines the amount of expansion of the star prior to the detonation, which determines the overall density profile (Jackson et al. 2010).

3.5. Mesh Geometry and Refinement

We performed our calculations in 2D $z - r$ cylindrical coordinates, extending radially from 0 to 65,536 km and along the axis of symmetry from −65,536 to 65,536 km. We selected a maximum refinement level corresponding to 4 km resolution using the PARAMESH adaptive mesh refinement scheme described in Fryxell et al. (2000). This resolution permitted efficiency in performing many repeated simulations with different initial conditions. The 4 km resolution was also informed by previous resolution studies in Townsley et al. (2007, 2009), which found that in 2D DDT simulations of C–O WD explosions the trends with resolution of total mass above the DDT density threshold at the DDT time are fairly robust. The amount of high-density mass at the DDT is important because it reflects the extent of neutronization during the deflagration and thus correlates with the IGE yield of the explosion. In addition, we wish to compare the hybrid IGE yields and other explosion characteristics with those of explosions from C–O WD progenitors previously explored in Krueger et al. (2012), which used 4 km resolution in FLASH with the same $z - r$ geometry we describe above. We can thus make our comparison robust by controlling for resolution and geometry factors.

4. SIMULATIONS AND RESULTS

4.1. Reference C–O WD Model

To compare the resulting features of SNe Ia produced by the hybrid model to those of previous studies of centrally ignited C–O WDs in the DDT paradigm (Krueger et al. 2010, 2012), we generate a reference C–O model. The reference model has the same central density as the hybrid model with a central temperature of $7 \times 10^8$ K, a core composition of $(^{12}$C = 0.4, $^{20}$Ne = 0.03, $^{16}$O = 0.57), and an envelope composition of $(^{12}$C = 0.5, $^{20}$Ne = 0.02, $^{16}$O = 0.48). For comparison, Figures 8 and 9 show the density and temperature profiles of the C–O and C–O–Ne hybrid WD.

We use the deflagration initialization of Krueger et al. (2012), which followed the method of Townsley et al. (2009). Townsley et al. (2009) reported that the scatter in $^{56}$Ni yield from SNe Ia could be produced in a theoretical sample of SN Ia simulations from C–O WD progenitors in the DDT scenario if the evolution of the flame surface shortly following the start of deflagration is described by a deformed spherical burned region within the progenitor core. This deformation consists of a perturbation of the sphere’s surface by a set of spherical harmonic functions with randomly chosen amplitudes. The theoretical SN Ia sample is formed by thus constructing a set of randomized burned geometries and then using each as the initial condition for a DDT simulation of the explosion.

We adopt this procedure to burn a spherical ignition region extending to a nominal radius of 150 km before applying a set of 35 randomly seeded amplitude perturbations, including those of Krueger et al. (2012), to ignite a suite of 35 C–O WD explosions. Finally, previous work has made the point that the explosion yield from simulations such as these may be sensitive to the initial conditions, motivating us to perform suites of 2D simulations instead of a modest number of 3D simulations (Calder et al. 2011). Although this kind of “vigorous” ignition is disfavored by convection zone ignition simulations (Nonaka et al. 2012), it appears necessary for DDT simulations to reproduce observed SNe (e.g., Seitenzahl et al. 2013; Sim et al. 2013).

4.2. Ignition Conditions for the C–O–Ne Hybrid WD

Given the temperature and composition profile of the hybrid model, if it is to undergo thermonuclear runaway, $^{12}$C ignition will begin at the base of the convective zone, where temperatures are highest. This coincides with a $^{12}$C abundance of $\approx 0.14$ at a stellar radius of 350 km. Because the hybrid progenitor model was produced by the one-dimensional MESA code and has only radial dependence, a shell-like ignition region naturally results when mapping the model to a multidimensional grid. Because this scheme cannot adequately
capture the effects of any low-mode, large-scale convection present in such a star, it is not clear what is the correct choice of ignition geometry for this hybrid WD model.

Lacking constraints on the exact geometry of the ignition region, we modify the ignition scheme described in Section 4.1 minimally for the hybrid progenitor by replacing the burned sphere with a thin burned shell modulated by a spherical harmonic function with variable amplitude and harmonic number, as shown in Figure 10 for different choices of the harmonic number. With this parameterization, the harmonic number controls the spatial localization and the amplitude controls the amount of initially burned mass. We therefore use our ignition parameterization to generate a suite of 35 hybrid realizations corresponding to a range of initially burned masses from \(3.0 \times 10^{-3} M_\odot\) to \(1.3 \times 10^{-2} M_\odot\) and a number of initially burned regions from 1 to 10. We note that our implementation of the initial conditions produced a slight asymmetry in the initially burned regions as a function of angle, as may be observed in Figure 10.

We demonstrate the influence of the number-and-amplitude parameters using the final \(^{56}\text{Ni}\) yield as a proxy for the explosion results in Figures 11 and 12. Because we use the same ignition geometry and conditions for the C–O realizations as were used in Krueger et al. (2012), the general size of the initially burned region for the C–O realizations remains very near \(0.0084 M_\odot\). However, because the nature of the ignition in the C–O–Ne hybrid progenitor is unknown, we chose to sample the number-and-amplitude parameter space to provide a range of initially burned masses for comparison. In spite of the scatter in Figure 11, we performed a linear fit between the \(^{56}\text{Ni}\) yield and initially burned mass for the C–O–Ne realizations, obtaining a slope indistinguishable from zero within uncertainties and an intercept that matches the average \(^{56}\text{Ni}\) yield for the C–O–Ne realizations (Table 1). We also show in Figure 12 that most of the variation in \(^{56}\text{Ni}\) yield from the C–O–Ne realizations originates from the interplay between the number of ignition regions and their size. The more ignition regions

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**Figure 10.** Distribution of initially burned mass for different choices of the harmonic number controlling the angular sinusoidal function. From left to right, the number of initially burned regions is 1, 2, 6, and 10. The harmonic number parameter is twice the number of ignition points. White denotes unburned fuel (\(^{12}\text{C}, ^{16}\text{O}, \) and \(^{20}\text{Ne}\)), and black denotes material in NSE (IGEs and \(\alpha\)-particles). The burned mass is located at a radial distance of 350 km, consistent with the location of maximum temperature at the base of the convective zone in the MESA profile used to initialize the FLASH simulations. For these four initial configurations, the average initially burned mass is \((6.67 \pm 0.02) \times 10^{-3} M_\odot\).

**Figure 11.** Dependence of the final \(^{56}\text{Ni}\) yield on the initially burned mass at ignition for the 35 C–O (red) and 35 C–O–Ne realizations (green).

**Figure 12.** Dependence of the final \(^{56}\text{Ni}\) yield on the distribution of the initially burned mass at ignition for the 35 C–O–Ne realizations.
that are used, the greater effect the variation on their size has on the spread in $^{56}$Ni yields.

Using these parameters to vary the initial burned mass and its distribution within the progenitor, we evaluate the effect of this parameterization on the estimated $^{56}$Ni yield, IGE yield, binding energy, and other explosion properties in the following section.

4.3. Characteristics of Explosions from C–O and C–O–Ne WD Models

We simulate the explosions of the hybrid and C–O realizations through the end of the detonation phase and compare their features in Figures 13, 14, and 16–20 below. We compare the $^{56}$Ni yields in the C–O and hybrid models, estimated from $Y_e$ and the NSE progress variable, by assuming that the composition in NSE is $^{56}$Ni plus equal parts $^{54}$Fe and $^{58}$Ni, as described in Townsley et al. (2009) and Meakin et al. (2009).

Production of $^{56}$Ni is comparable between the C–O and hybrid cases (Figure 13), with the full range of values from each suite of simulations shown in the shaded regions and the mean values shown by solid curves. The DDT event can be distinguished in the $^{56}$Ni evolution by the sharp increase in the rate of $^{56}$Ni production around 1.5 s that rapidly yields over 0.5 $M_\odot$ of $^{56}$Ni. While the C–O cases show a wider variation in the time at which the DDT occurs, these also have a narrower spread in final $^{56}$Ni mass relative to the hybrid models. The hybrid models also tend to produce more $^{56}$Ni in the deflagration phase, and some of them show a temporary plateau in $^{56}$Ni production between 1.5 and 2 s. The same feature is also evident in the binding energy curves of Figure 14, computed by summing the realization’s gravitational potential, internal, and kinetic energies.

This feature is a peculiarity of the off-center ignition in the hybrid models that is absent in the C–O cases and results from the relatively $^{12}$C-poor, cooler core region burning about 0.25 s after the detonation front has swept through the rest of the star. This delayed burning is shown in Figure 15, which demonstrates the progression of the detonation front into the core. Although a feature evolving over so short a time this early in the explosion will likely not be visible in the SN light curves, the delayed contribution of the core to $^{56}$Ni production may modify the $^{56}$Ni distribution in space and velocity, potentially yielding spectral differences compared to nondelayed hybrid as well as C–O WD explosions.

The dynamical qualities of the explosion shown in the binding energy curves of Figure 14 indicate that the time distribution of unbinding is narrower for the hybrid models than for the C–O models, though the hybrid models have a wider distribution of final binding energies in all cases lower than the binding energies of the C–O models within 1 s of becoming unbound. This should correlate to a lower expansion velocity of the ejecta and thus slower cooling and delayed transparency relative to ejecta from C–O models. The binding energy curves also explain the differences in expansion the models undergo during deflagration and detonation, shown by the mass above the density threshold $2 \times 10^7$ g cm$^{-3}$ in Figure 16. For times prior to $\approx 1.2$ s, the C–O mass curves lie slightly lower than the average hybrid mass curve, indicative of a greater degree of expansion on average for the C–O models.

However, the hybrid mass curve range encompasses that of the C–O mass curves until $\approx 1.4$ s, reflective of the fact that until then, some hybrid realizations are more tightly bound than all the C–O realizations due to burning less mass and thus expanding less. During the detonation phase, however, the C–O models show a much wider variation in expansion than do the hybrid models in spite of having a smaller range of kinetic energies and mass burned to IGEs (Figures 14 and 20) once unbound. This is due to the C–O models demonstrating a much wider range of DDT times than the hybrid models.

| Progenitor Type | $^{56}$Ni ($M_\odot$) | IGE ($M_\odot$) | Kinetic Energy ($\times 10^{51}$ ergs) |
|----------------|---------------------|----------------|-------------------------------------|
| C–O            | 0.97 ± 0.06         | 1.12 ± 0.07    | 1.39 ± 0.05                         |
| C–O–Ne         | 0.86 ± 0.10         | 0.98 ± 0.11    | 1.06 ± 0.10                         |
Figure 15. Progress of the burning front into the stellar core for one hybrid C–O–Ne realization, delayed relative to complete burning throughout the rest of the star. For reference, the initially burned geometry is shown at left. Material is shaded based on the reaction progress variables so that white denotes unburned fuel ($^{12}$C, $^{16}$O, and $^{20}$Ne) and red denotes ash from $^{12}$C and $^{20}$Ne burning. Green then denotes material in NSQE (primarily intermediate-mass silicon-group elements), and black denotes material in NSE (IGEs and $\alpha$-particles). From left to right, the burning is shown at 0.0, 1.9, 2.0, and 2.1 s. The DDT time for this realization is $\approx 1.4$ s.

Figure 16. Evolution of the total mass having density greater than $2 \times 10^7$ g cm$^{-3}$ for C–O and hybrid C–O–Ne WD realizations. The time-averaged value among C–O realizations is shown in yellow, with the full range of values at any point in time for the C–O realizations shown in red. Likewise, the time-averaged value among C–O–Ne realizations is shown in blue and their range of values shown in green.

Figure 17 compares the final IGE yield of the C–O and C–O–Ne models with the degree by which the models expand during the deflagration phase. The latter is characterized by the mass above $2 \times 10^7$ g cm$^{-3}$ at the DDT time, with more high-density mass indicating less expansion during deflagration. The averages of both the C–O and C–O–Ne suites along both axes are indicated by the shaded regions with ±1σ widths. The trend for both C–O and C–O–Ne models is that less expansion during the deflagration phase results in greater IGE yields, expected because low expansion results in there being more high-density fuel for the detonation to consume. In addition, both the C–O and C–O–Ne models expand over similar ranges during deflagration on average, showing that they are dynamically comparable in spite of having qualitatively different deflagration ignition geometries. Furthermore, for similar deflagration expansion, the C–O models tend to yield consistently greater IGE mass, suggesting that the lower IGE yields from C–O–Ne models are not a result of these models expanding differently than the C–O models. Rather, we interpret this disparity as indicating that the lower IGE yield in C–O–Ne models results from their lower $^{12}$C abundance and the fact that given similar fuel density, their $^{20}$Ne-rich fuel will burn to cooler temperatures than fuel in the C–O models. This in turn will result in slower burning to IGEs and thus a lower IGE yield.

The estimated $^{56}$Ni yields are shown in Figure 18 across the range of masses burned to IGEs for all C–O and C–O–Ne realizations at 4.0 s simulation time, at which point the total mass burned to $^{56}$Ni had become constant (see Figure 13). For comparable masses burned to IGEs, the hybrid models tend to
Figure 18. Production of $^{56}$Ni and mass burned to IGEs for C–O (red) and hybrid C–O–Ne (green) WD realizations. To estimate the overall fraction for each case, a linear fit is shown for C–O (yellow) and C–O–Ne (blue) realizations.

Figure 19. Estimated fraction by mass of IGE material producing $^{56}$Ni evolving in time for C–O and hybrid C–O–Ne WD realizations. The time-averaged value among C–O realizations is shown in yellow, with the full range of values at any point in time for the C–O realizations shown in red. Likewise, the time-averaged value among C–O–Ne realizations is shown in blue and their range of values is shown in green.

Figure 20. Mass burned to IGEs for C–O and hybrid C–O–Ne WD realizations evolving in time. The time-averaged value among C–O realizations is shown in yellow, with the full range of values at any point in time for the C–O realizations shown in red. Likewise, the time-averaged value among C–O–Ne realizations is shown in blue and their range of values shown in green.

consistently produce slightly more $^{56}$Ni than the C–O models, although the ratio of IGE mass producing $^{56}$Ni given by the slope is the same in both cases, within the fit error. The reason for this trend is evident from Figure 19, which shows the fraction by mass of IGE material producing $^{56}$Ni evolving in time, and Figure 20, which shows the concurrent evolution of mass burned to IGE. During the deflagration phase, the C–O–Ne models on average burn more material to IGEs and also had a significantly higher fraction of IGE material producing $^{56}$Ni, yielding more $^{56}$Ni than the C–O models. This may be due to greater neutronization in the early deflagration of the C–O models, which are ignited closer to the center and thus at slightly higher density than the initial deflagration of the C–O–Ne models. However, during the subsequent detonation phase, the C–O models on average burn more mass to IGE while maintaining a $^{56}$Ni/IGE fraction very similar to that of the C–O–Ne models, yielding significantly more $^{56}$Ni by the end of the detonation phase. For reference, Table 1 summarizes the average values of $^{56}$Ni yield, IGE yield, and final kinetic energy for the 35 C–O and 35 C–O–Ne realizations with one standard deviation uncertainties.

5. CONCLUSIONS

Our simulations of thermonuclear (Type Ia) SNe from both hybrid C–O–Ne and reference C–O WD progenitors using the DDT paradigm have shown that on average the hybrid progenitors yield 0.1 $M_{\odot}$ less $^{56}$Ni than the C–O WDs. While this indicates that SNe Ia from C–O–Ne hybrids will be dimmer on average than those from C–O WDs, we also find sufficient variance in burning efficiency with the geometry of the ignition region precipitating thermonuclear runaway such that there are some hybrid progenitors that yield more $^{56}$Ni than some C–O progenitors. Furthermore, we have found that not only do hybrid C–O–Ne progenitors deposit an average of 24% less kinetic energy in their ejecta than C–O progenitors but also this trend of more weakly expelled ejecta from hybrids is robust across all ignition geometries. The consistency of this result suggests that it is a consequence of the lower energy release from Ne burning compared to C burning in spite of the fact that using $^{20}$Ne as an alternate fuel can still yield comparable $^{56}$Ni production in some cases.

As we noted above, we found considerable variation in the $^{56}$Ni production for both hybrid and traditional C–O models, and in particular, we found a much wider range of DDT times in the C–O models than the hybrid models. While in the realm of speculation, this result could follow from a greater degree of randomization in the geometry of the initially burned region for the C–O models than in the hybrids. The C–O models are initialized with an amplitude perturbation of the initially burned region composed of several angular modes, whereas the thickness of the initially burned region in each of the hybrid models is controlled by a single angular mode.

We also found that for some ignition geometries in the hybrid progenitor, a combination of off-center ignition, flame buoyancy, and composition permits their cooler core region to
delay burning until nearly 0.25 s after the detonation front has consumed the rest of the star. This result is unique to the hybrid progenitors that rely on the Urca process to provide a lower bound on the convective zone outside the core of the WD. Delayed core burning in these WDs may result in a modified $^{56}\text{Ni}$ distribution in their ejecta compared to ejecta from C–O WDs or even other hybrids with prompt core burning. Exploration of such effects is the subject of future work.

As our explosions from hybrid progenitors have a lower $^{56}\text{Ni}$ yield and hence lower brightness than traditional C–O models, the question of these events as the source of observed subluminous events, e.g., Type Iax SNe (Foley et al. 2013), arises. Our finding of an average $^{56}\text{Ni}$ yield of 0.1 $M_\odot$ less than the C–O (and the larger range of yields) indicates that explosions from these progenitors in the DDT paradigm are not subluminous and cannot on their own explain subluminous events like Type Iax SNe.

A recent study by Kromer et al. (2015) addressed pure deflagrations in near-Chandrasekhar-mass hybrid WDs as the possible progenitor systems of these faint events. The study found that most of the mass stays bound and that early epoch light curves and spectra calculated from the explosion models are consistent with observations of SN 2008ha (Foley et al. 2009). We note that comparison between our results and these is difficult for reasons besides the obvious difference of the detonation phase in our simulations. The near-Chandrasekhar-mass progenitor model of Kromer et al. (2015) is substantially different in that it is parameterized and it does not include the effects of late-time convection, or the Urca process. Also, the ignition of the deflagration is substantially different. For these reasons, there is limited utility in a direct comparison between results.

Another recent study by Bravo et al. (2016) investigated explosions from two classes of hybrid WDs, those with relatively large central C–O regions that follow from off-center C burning that stops prior to consuming the core due to neutrino cooling and those with somewhat smaller C–O cores that follow from CBM inhibiting burning. The latter “medium-sized” cores are most similar to our models as they assume similar stellar evolution.

The group performed suites of one-dimensional simulations evoking explosions via both pure detonations and DDTs and including “homogeneous” models that include the effect of mixing during C simmering prior to the ignition of the explosion. Some of the DDT models produce $^{56}\text{Ni}$ yields consistent with our results, but the model closest to ours (number 28) produces a yield substantially lower than our average result. This difference follows from the choice of $\rho_{\text{DDT}}$ and its very different role in one dimension compared to two dimensions. DDT in two dimensions occurs at the tops of plumes, so that the degree of pre-expansion for a particular $\rho_{\text{DDT}}$ is significantly less in two dimensions than in one dimension. Thus, 2D simulations produce significantly higher $^{56}\text{Ni}$ yields than one-dimensional simulations for a similar $\rho_{\text{DDT}}$. For similar reasons, one-dimensional simulations with a higher $\rho_{\text{DDT}}$ give comparable yields to 2D simulations with lower $\rho_{\text{DDT}}$, leading to some cases with comparable yields.

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