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TYPE II SUPERNOVA

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Theoretical Models for Type I and Type II Supernova

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Abstract

Recent theoretical progress in understanding the origin and nature of Type I and Type II supernovae is discussed. New Type II presupernova models characterized by a variety of iron core masses at the time of collapse are presented and the sensitivity to the reaction rate $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$ explained. Stars heavier than about $20 M_{\odot}$ must explode by a "delayed" mechanism not directly related to the hydrodynamical core bounce and a subset is likely to leave black hole remnants. The isotopic nucleosynthesis expected from these massive stellar explosions is in striking agreement with the sun. Type I supernovae result when an accreting white dwarf undergoes a thermonuclear explosion. The critical role of the velocity of the deflagration front in determining the light curve, spectrum, and, especially, isotopic nucleosynthesis in these models is explored.

1. New Presupernova Models of Massive Stars

As a part of an ongoing project to calculate and understand the final evolution of massive stars and the consequent Type II supernova display, we have recently simulated the evolution of stars of a variety of masses: 11, 12, 15, 20, 25, 50, and $100 M_{\odot}$ (Table 1). Some aspects of these models have been published previously (Wilson et al. 1985; Weaver and Woosley 1985) and the detailed evolution of the remainder is the subject of papers now in preparation. Here we briefly review the properties of the presupernova models and, especially the sensitivity of structure and nucleosynthesis to the new reaction rate for $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$. Stars have been studied using a revised version of the stellar evolution program KEPLER (Weaver, Zimmerman, and Woosley 1978; Weaver, Woosley, and Fuller 1984) which incorporates the new weak interaction rates of Fuller, Fowler, and Newman (1982ab, 1985), an improved implementation of weak processes at an earlier stage of the star's life (near oxygen core depletion), revisions to nuclear screening corrections (although see paper by Mochkovitch, this volume; photodisintegration reactions should not be screened), finer time step and

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zoning criteria, and Cameron (1982) initial abundances. Because of the smaller helium abundance employed in the new studies, helium core masses are smaller than in some of our earlier studies. For example, the new 11 $M_{\odot}$ star has a helium core at the end of its life of 2.4 $M_{\odot}$, slightly smaller than the helium core in our old 10 $M_{\odot}$ star (Woosley, Weaver, and Taam 1980). Helium core masses in the larger stars are also reduced by about 10%. More importantly, the revised weak rates and prescription for their implementation leads to iron core masses that are considerably smaller in 12 and 15 $M_{\odot}$ stars than in Weaver, Zimmerman, and Woosley (1978). Our present 15 $M_{\odot}$ model (Table 1) is an example with an iron core mass reduced from its 1978 value, 1.56 $M_{\odot}$, to 1.33 $M_{\odot}$. On the other hand, revisions to the reaction rate for $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$ lead to larger iron core masses in the bigger stars. The new 25 $M_{\odot}$ model, for example, which had a mass in 1978 of 1.61 $M_{\odot}$ and in Weaver, Woosley, and Fuller (1985) of 1.35 $M_{\odot}$, now has a mass of 2.1 $M_{\odot}$. These differences will be explored later. First we deal with the evolution of the lighter stars.

TABLE 1
PRESUPERNOVA MODELS AND EXPLOSIONS

Main Sequence Mass ($M_{\odot}$)	Iron Core Remnant ($M_{\odot}$)	Explosion Energy ^a (10^{50} erg)	Baryon Mass ^a ($M_{\odot}$)	Neutron Star ^a ($M_{\odot}$)
11	— ^b	3.0	1.42	1.31
12	1.31	3.8	1.35	1.26
15	1.33	2.0	1.42	1.31
20	1.80	—	—	—
25	2.10	4.0	2.44	1.96
50	2.45	—	—	—
100	— ^b	~4	39	Black Hole

^a All except for 100 $M_{\odot}$ determined by Wilson et al. (1985).

^b Never developed iron cores in hydrostatic equilibrium.

Our new 11 $M_{\odot}$ star had an evolution that differed markedly from all the other stars examined here. It was also the only model reported here which employed an old (Fowler, Caughlan, and Zimmerman 1975) rate for the $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$ reaction rate, although, in this particular case, we do not view that as critical. Qualitatively, the life history was similar to that which has been reported previously for a 9.6 $M_{\odot}$ star by Nomoto (1984ab; see also Miyaji et al. 1980). The helium core mass in the Nomoto study was 2.4 $M_{\odot}$, roughly the same as here at the end of the evolution. However, just after depletion of carbon in the center of our 11 $M_{\odot}$ star, the helium core mass

was $2.6 M_{\odot}$, hydrogen dredge-up of the helium core during off-center carbon burning accounts for the difference. The later stages of evolution are characterized by temperature inversion in the core (owing especially to plasma neutrino losses) and off-center burning. A neon core develops (60% neon and 20% each of magnesium and oxygen) containing very nearly a Chandrasekhar mass ($1.45 M_{\odot}$). That this mass exceeds $1.37 M_{\odot}$ marks a critical departure between the present work and that of Nomoto (1984ab). For the smaller mass, neon and subsequent burning stages never ignite (under stable conditions) whereas in our $11 M_{\odot}$ model, both neon and oxygen burning do occur in a shell from about $0.6 M_{\odot}$ to $1.3 M_{\odot}$ (Fig. 1) and, in a second convective burning stage from $1.3 M_{\odot}$ to $1.4 M_{\odot}$. Thus at the end of its life this star consists of a cold ($T \sim 5 \times 10^8$ K) central core of about $0.5 M_{\odot}$ of neon surrounded by silicon and sulfur with a total mass slightly exceeding the (zero entropy) Chandrasekhar mass. Further cooling then leads to the collapse of the core containing combustible nuclear fuel.

It must be acknowledged at this point that the onset of collapse in this star has not been calculated in an entirely self-consistent fashion. Electron capture reactions were not included in the calculation of neon and oxygen burning although oxygen burns at such a high density off-center here that such captures would certainly have been important, probably triggering core collapse at an earlier stage. At a density of $\sim 10^8$ g cm $^{-3}$, for example, capture on the products of oxygen burning (e.g. $^{33}\text{S}(e^-, \nu)^{33}\text{P}$; see Woosley, Arnett, and Clayton 1972) would lead to $\langle Y_e \rangle \sim 0.49$ reducing the zero entropy Chandrasekhar mass to $1.38 M_{\odot}$ and causing the core to collapse. We do not think that the continued evolution is especially sensitive to this detail, but a self-consistent calculation would obviously be desirable. It is interesting to note that the contraction of the core is not without nucleosynthetic import in the overlying layers. We find that the heating of the helium burning shell leads to extension of the helium convective zone into the hydrogen envelope. The mixing of protons into a super-heated helium shell has obvious and far reaching implications for r and s-process nucleosynthesis which will be discussed elsewhere.

As the core contracts to 2.5×10^{10} g cm $^{-3}$, oxygen burning, greatly enhanced by electron screening, ignites at a temperature of $\sim 5 \times 10^8$ K. We note that at this point the core is also hovering on the verge of general relativistic instability (Shapiro and Teukolsky 1983) because Γ is so nearly equal to $4/3$. Future calculations should include post-Newtonian gravity. Owing to the extremely degenerate nature of the core the nuclear runaway that ensues leads to complete combustion to iron group nuclei. Thereafter the evolution becomes quite similar to that described by Nomoto (1984ab) for his $2.4 M_{\odot}$ helium core. Iron group products capture electrons efficiently leading to a sudden drop in Y_e at the center of the star to about 0.40. The thermal increment to the pressure from burning is inconsequential compared to the great loss in electron degeneracy pressure and the core begins to implode dynamically. As neon and magnesium and, later, silicon and sulfur, fall down they are heated by compression

Fig. 1 - Composition (a) and structure (b) of an $11 M_{\odot}$ star just after the off-center ignition of oxygen burning. The temperature at the base of the convective oxygen shell at this point is 1.87×10^9 K. The central density and temperature are $8.19 \times 10^7 \text{ g cm}^{-3}$ and 5.63×10^8 K respectively. The surface luminosity, effective temperature, and radius, all of which should not change prior to the stellar explosion, are $1.47 \times 10^{38} \text{ erg s}^{-1}$, 4290 K, and $2.47 \times 10^{13} \text{ cm}$. The neutrino luminosity is $1.21 \times 10^{42} \text{ erg s}^{-1}$. Note scale breaks at $3.0 M_{\odot}$ (a) and $4.0 M_{\odot}$ (b).

and burn explosively so that a standing combustion front exists at ~ 110 to 170 km. Once the collapse is well under way, energy transport by convection is negligible.

The explosion of this model was calculated by Wilson et al. (1985) and the reader should see that reference for details. We note, however, that the explosion was of the "delayed" variety, occurring only after about 0.07 s. This differs qualitatively from the results of Hillebrandt, Nomoto, and Wolff (1984) who obtained a prompt hydrodynamical explosion for Nomoto's $9.6 M_{\odot}$ star. We suspect that the difference may lie in different nuclear equations of state employed in the two codes or perhaps slight differences in the initial models. This point is under investigation by Wilson and Mayle.

Other models were also calculated with the improved version of KEPLER but using the recently revised rate for $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$. Measurements by Kettner et al. (1982) and a reanalysis of new and old data by Langanke and Koonin (1984) have led to a revised rate, as tabulated by Caughlan et al. (1985), that under typical conditions in massive stars, is about 3 times the old value used in our previous calculations (Fowler, Caughlan, and Zimmerman 1975). Such a large revision has, as we shall see, major implications not only for nucleosynthesis, but also for the structure of the presupernova star. The iron core mass in the 12 and $15 M_{\odot}$ models are not greatly altered from those obtained (for 20 and $25 M_{\odot}$ stars) by Weaver, Woosley, and Fuller (1984), although, as it turns out, the extent of the oxygen burning shell is much larger with the new rate), but the new $25 M_{\odot}$ model has a much larger iron core (Figs. 2 and 3).

Why should there be such great variation in the iron core mass for stars of 15 and $25 M_{\odot}$ or for two stars, both of $25 M_{\odot}$, differing only in the magnitude of the reaction rate for $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$? The answer resides in the nature of carbon and neon burning and how they affect the entropy structure of the stellar core (Fig. 4). In the $15 M_{\odot}$ star and in the older $25 M_{\odot}$ star with low α -rate, carbon burning ignites as a well developed, exoergic, convective burning stage. In fact, the $15 M_{\odot}$ star goes through three distinct stages of carbon convective burning before igniting neon burning at its center. The first stage depletes carbon in a region out to about $0.4 M_{\odot}$; the second convective shell goes from ~ 0.4 out to $1.0 M_{\odot}$; and finally a third stage burns carbon out to about $1.5 M_{\odot}$. The initial carbon abundance following helium burning in this star is 0.14. While carbon burning goes on, neutrino losses cool the core and its entropy decreases. During the first carbon convective stage, the central conditions are $T_c = 8.1 \times 10^8$ K, $\rho_c = 3.8 \times 10^5$ g cm $^{-3}$ while at the end of the third stage, (at neon ignition) $T_c = 1.6 \times 10^9$ K, $\rho_c = 8.9 \times 10^6$ g cm $^{-3}$. This loss of entropy from the core allows the evolution to "keep in touch" with the Chandrasekhar mass (as adjusted for the finite entropy of the electrons). Thus a strong carbon burning shell is established at the edge of a (semi-)degenerate core. From this point onwards there always exists a sharp increase in the entropy at $M \sim 1.4$ to $1.5 M_{\odot}$ that restrains the outward extension of convective shells during oxygen and silicon burning so that when the

Fig. 2 - Composition (a) and structure (b) of a $15 M_{\odot}$ star at the onset of core collapse (as defined by collapse velocity equals 1000 km s^{-1}). Interior to the iron core ($1.33 M_{\odot}$) energy is being lost due to a combination of neutrino emission from electron capture and photodisintegration. Peaks in the nuclear energy generation, ϵ_{nuc} , are apparent at the silicon, oxygen, and helium burning shells. Also plotted are the neutrino losses due to plasma processes (ϵ_{ν} in the region outside the iron core). Note the rapid fall off in density just outside the iron core and the substantial abundances of the elements silicon thru calcium in the oxygen convective shell.

Fig. 3 - Composition (a) and structure (b) of a 25 $M_{\odot}$ star at the onset of core collapse. Note the large iron and the less rapid fall off in density (compared to Fig. 2) outside that iron core. Notation and sampling time are as in Fig. 3.

Fig. 4 - Entropy distributions in the inner regions of $15 M_{\odot}$ and $25 M_{\odot}$ model star just prior to carbon ignition (solid line), neon ignition (dashed line), and oxygen ignition (dot-dashed line). The total entropy and partial entropy in the electrons is given for both stars in dimensionless units of Boltzmann's constant. Stellar structure is most sensitive to the electronic entropy since the electrons are the major source of pressure. The locations of convective shells are sensitive to the distribution of total entropy. Entropy decreases substantially in the $15 M_{\odot}$ model during carbon burning.

star finally collapses it does so with an iron core of $1.33 M_{\odot}$, close to the traditional Chandrasekhar mass for material with equal numbers of neutrons and protons. At least a portion of the decrease from ~ 1.44 to $1.33 M_{\odot}$ is a consequence of electron capture during and shortly after oxygen burning.

The $25 M_{\odot}$ star with the most recent α -capture rate (Table 1) is quite another story. Because the carbon abundance following helium depletion is so low, $\sim 9\%$ by mass, carbon burning and neon burning never ignite in the center of the star as exoergic, convective burning stages. The trace abundances of carbon and neon burn away radiatively, without the nuclear energy generation ever exceeding neutrino losses. Because there is no cooling stage, and because the $25 M_{\odot}$ star had a larger entropy in its core to begin with, the core does not become especially degenerate (Fig. 4) and is not sensitive to the Chandrasekhar mass. Carbon is depleted radiatively out to a mass of about $2.5 M_{\odot}$. Neon burns radiatively out to about $1.5 M_{\odot}$. The next fully developed burning stage after helium burning is oxygen burning. Prior to silicon core ignition, oxygen burns in a convective core first out to about $1.35 M_{\odot}$ and then in a convective shell out to $2.4 M_{\odot}$. When silicon does ignite, it is within a $2.4 M_{\odot}$ core comprised of almost pure silicon and sulfur. The large entropy increase associated with the oxygen shell is then at 2.4 , not $1.4 M_{\odot}$ as it was in the $15 M_{\odot}$ model. Silicon burns out to $1.3 M_{\odot}$ but, even with $Y_e \sim 0.46$, the core is too small to collapse given its thermal content. Thus a silicon convective shell burns out to $2.1 M_{\odot}$ before the core collapses. It is interesting that the property of a massive star that most sensitively determines its final evolutionary state, neutron star or black hole, namely the iron core mass at collapse, is so sensitive to occurrences during its relatively early life.

The $50 M_{\odot}$ stellar evolution was qualitatively similar to that of the $25 M_{\odot}$ model although, owing to a larger entropy, less centrally condensed core, and smaller carbon abundance, the iron core at the end was even larger, $2.45 M_{\odot}$.

The $100 M_{\odot}$ star, on the other hand, met a death unlike any reported thus far. Near the end of helium burning, the hydrogen envelope of the star was removed from the calculation when it became apparent that no reasonable choice of surface boundary conditions would allow it to stay on with a luminosity that was very nearly super-Eddington. Thus in its final stages the star was a $42 M_{\odot}$ helium core, i.e., a massive Wolf-Rayet star. Removal of the envelope has no important effect on the subsequent evolution of the helium core although it obviously affects the observational properties of the star. Following the central (radiative) exhaustion of carbon and neon, this core encountered the electron-positron pair instability upon attempting to ignite oxygen. The first collapse resulted in a peak central temperature of 3.0×10^9 K at a density of $1.5 \times 10^{-6} \text{ g cm}^{-3}$ which led to explosive oxygen burning and expansion. The initial explosion was far too weak to unbind the entire star, but a portion, about $1/4 M_{\odot}$, was ejected from the surface at velocities of several thousand kilometers per second (2.2×10^{49} erg). The central density of the star declined to $7.1 \times 10^{-4} \text{ g cm}^{-3}$,

and roughly one month long Kelvin-Helmholtz stage ensued as the star contracted and encountered the pair instability again. This time a higher peak temperature, 3.3×10^9 K is reached and a more violent explosion ensues, 6.5×10^{49} erg carried by an additional quarter solar mass of ejecta. For a time, again of about a month, the star oscillates violently but eventually settles down to encounter the pair instability a third time, the peak temperature this time being 3.7×10^9 K. This leads to a strong explosion and the ejection of $2.7 M_{\odot}$ of surface material (helium, carbon, and oxygen) with energy 3.4×10^{50} erg. Oscillations and relaxation then lead to a fourth instability with $T = 4.6 \times 10^9$ K and a weak explosion. On the fifth time down, 2 years after the onset of the first explosion, there is nothing left at the center of the star to burn, hence no explosion. The remainder of the star, neglecting rotation, becomes a black hole. This model and its observational consequences should amuse and employ us theoreticians for years. Stay tuned!

As a general result of our study of stars in the 11 to $100 M_{\odot}$ range, it appears that the compact remnants of stars having main sequence mass $\gtrsim 20 M_{\odot}$ [with exact value sensitive to $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$], will have quite different properties from those of lower mass. In particular, since nuclear equations of state suggest an upper bound to the mass of a stable neutron star of $\sim 2.0 M_{\odot}$, stars heavier than $\sim 20 M_{\odot}$ may leave black holes and lighter stars (but heavier than $\sim 10 M_{\odot}$) will leave neutron stars. It is important also to note that an object which eventually, after cooling and deleptonization, becomes a black hole may be the residual of an explosion which ejects matter, produces explosive nucleosynthesis, and exhibits a light curve not markedly discrepant with observations of Type II supernovae (Wilson et al. 1985).

2. Delayed Explosions of Type II Supernovae

For many years the outward propagation of the shock wave produced by the bounce of these iron cores has been studied as a possible mechanism for the explosion of the star (cf. Bowers and Wilson 1982; Arnett 1980, 1983; Brown, Bethe, and Baym 1982; Hillebrandt 1984; Bruenn 1984; Cooperstein 1982; Kahana, Baron, and Cooperstein 1984). For the most part, the results of these studies have not been particularly encouraging, except, perhaps, in the case of very low mass iron cores (Hillebrandt 1982; Hillebrandt, Nomoto, and Wolff 1984; Cooperstein 1982). The shock stalls, overwhelmed by photodisintegration and neutrino losses, and the star does not explode. More recently, slow late time heating of the envelope of the incipient neutron star has been found to be capable of rejuvenating the stalled shock and producing an explosion after all (Wilson 1984; Bethe and Wilson 1985; Wilson et al. 1985). The basic mechanism is energy transport by electron neutrinos and anti-neutrinos which capture on nucleons and scatter on electron-positron pairs just behind the stalled (accretion) shock. The source of the neutrinos is the cooling proto-neutron-star beneath the

shock and, to a lesser but important extent, the neutrinos liberated by the accretion flux itself. A complete discussion of the complexities (and attendant uncertainties) associated with this particular explosion mechanism are best left to the referenced papers by Wilson and co-workers and to future papers. For now we merely note that each of the explosions listed in Table 1 were calculated by Wilson and Mayle and were of this "delayed" variety (Wilson et al. 1985). Credible variations in the nuclear equation of state may allow the prompt explosion of the 11, 12, and 15 $M_{\odot}$ models (E. Baron and H. Bethe; private communication; Bethe and Brown 1985), but it appears highly unlikely that explosion of iron cores as large as 2.1 $M_{\odot}$ (i.e. the 25 $M_{\odot}$ model) will ever be achieved solely by a hydrodynamical bounce. Thus the nucleosynthesis we shall now discuss is quite contingent upon the successful occurrence of a delayed explosion. Otherwise the complexity of rotation will have to be considered (cf., Bodenheimer and Woosley 1983).

3. Nucleosynthesis in Type II Supernovae

The passage of the shock wave through the overlying mantle of the star, in addition to providing the impulse for its ejection, leads to high temperatures and nuclear reactions that have been followed in detail for the 15 and 25 $M_{\odot}$ models (Weaver and Woosley 1985; Wilson et al. 1985). Figure 5 shows the explosive nucleosynthesis in the 25 $M_{\odot}$ model and nucleosynthetic results of 15 and 25 $M_{\odot}$ models are both shown compared to solar abundances in Figure 6. In the case of the 15 $M_{\odot}$ model, the abundances ejected are almost entirely, with the obvious exception of the iron group, produced in the pre-explosive stages of evolution and merely shoved off the star by the explosion. The diminished importance of explosive nucleosynthesis, as compared, for example, to past studies (Weaver and Woosley 1980) is a consequence both of the low explosion energy and the rapid fall off in density around the core of the preexplosive star. The large abundances of silicon through calcium are especially a result of an extensive oxygen burning convective shell just outside the core (Fig. 2). It is interesting to note that the energy liberated by nuclear reactions was not an inconsequential fraction of the final explosion energy in the 15 $M_{\odot}$ model. The final kinetic energy at infinity was calculated to be 2×10^{50} erg, 30% of which was generated by nuclear burning during shock wave passage.

Because of the larger explosion energy in the 25 $M_{\odot}$ model, there was a considerable amount of explosive nucleosynthesis and nuclear energy generation formed an even greater fraction of the final kinetic energy at infinity. The initial shock contained roughly 10^{51} erg as calculated in the core bounce program, but the binding energy of the mantle was also about 10^{51} erg. Nuclear burning gave 4×10^{50} erg and the final kinetic energy was also 4×10^{50} erg. Thus without the energy from nuclear burning following shock wave passage a portion of this star may well have reimploded. In the

Fig. 6 - Comparison of bulk elemental nucleosynthesis in 15 and 25 $M_{\odot}$ supernovae to solar abundances (Cameron 1982). The enhancement factor is the mass fraction of a given element in the ejecta compared to its mass fraction in the sun (which was also its initial abundance in the calculation). The deficient production of carbon, neon, and magnesium, especially in the 15 $M_{\odot}$ study, is a consequence of the revised rate for $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$. The large productions of silicon thru calcium in the 15 $M_{\odot}$ model do not exist if this material falls back onto the neutron star following hydrodynamical interaction with the stellar envelope. Such was the result in the present study, but this occurrence is quite model dependent.

Fig. 5 - Composition of the inner regions of a $25 M_{\odot}$ supernova following shock passage and ejection. The pre-explosive composition (see Fig. 3) is substantially altered out to about $3.6 M_{\odot}$. Composition interior to the mass cut, $2.44 M_{\odot}$, is not plotted. The explosion was calculated by Wilson et al. (1985).

15 $M_{\odot}$ model calculated by Wilson et al. (1985) some mass reimplosion was actually observed.

Figure 7 shows the isotopic nucleosynthesis resulting from the explosion, by Wilson's delayed mechanism (Wilson et al 1985), of our 25 $M_{\odot}$ presupernova model. The comparison with solar abundances is very good, much better, for example, than that published for a previous 25 $M_{\odot}$ model by Woosley and Weaver (1982). The changes reflect principally the altered structure of the presupernova mantle and core brought about by alterations in weak interaction rates (Weaver, Woosley, and Fuller 1984) and in the reaction rate for $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$. It is worth noting that 34 species out of a total of 61 in this mass range are produced within a factor of two of their relative solar abundances and 52 are produced within a factor of 4. In the sun these abundances span a range in mass fraction of 7 orders of magnitude. Of the remaining 9, ^{13}C and ^{15}N are probably the products of hydrogen burning in lower mass red giant stars and novae respectively (cf. Wallerstein 1973; Lambert and Tomkin 1974; Sneed and Lambert 1975; Wallace and Woosley 1981); the origin of fluorine is unknown (although see Howard, Arnett, and Clayton 1971; Scalo and Despain 1976); ^{18}O is produced in a 15 $M_{\odot}$ star to be reported elsewhere; ^{46}Ca and ^{47}Ti are very rare species whose production might be quite sensitive to poorly determined nuclear reaction rates (or a change in stellar mass); ^{48}Ca and ^{54}Cr can be made in a neutron-rich nuclear statistical equilibrium (Hainebach et al. 1974; Hartmann, Woosley, and El Eid 1985); and ^{63}Ni was near the end of our nuclear reaction network and may not have been tracked accurately.

Although the agreement is really very good overall, some annoying discrepancies can be noted. Carbon is down relative to oxygen. That's not so bad. Perhaps carbon is produced in low mass stars (Dufour 1984). Of greater concern are the products of carbon burning, eg. ^{20}Ne , ^{23}Na , and ^{24}Mg . It appears unlikely that all these abundances could systematically be underestimated by a factor of two in the sun (relative to oxygen, say), nor does it appear probable that other masses of supernovae, eg. 15 $M_{\odot}$ are going to fill in the gap. All things considered, it would have been better if the rate for $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$ had not increased quite so much. Further increases (beyond the roughly factor of 3 times FCZ75 used here) would make matters worse still. We shall see. The rates will, of course, be whatever the nuclear physicists determine. It is interesting however, that in our previous model (Weaver, Woosley, and Fuller 1984), these same species were major overproductions while the silicon through calcium group was underproduced; the only difference, a smaller rate for $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$.

Following the observations of Mahoney et al. (1982), increased attention has been focused upon the production of ^{26}Al in Wolf Rayet stellar winds (see papers by Casse and by Dearborn this volume), novae (see paper by Truran this volume and Wallace and Woosley 1981) and supernovae. Clayton (1984), in particular, has emphasized the difficulty associated with producing the (strong) observed signal using ^{26}Al

Fig. 7 - Isotopic nucleosynthesis in a 25 $M_{\odot}$ explosion. Final abundances in the ejecta are plotted for isotopes from ^{12}C to ^{64}Ni compared to their abundances in the sun (Cameron 1982). An average overproduction factor of 9 characterizes the distribution. If one gram in 9 of the matter in the Galaxy has experienced conditions like those in a 25 $M_{\odot}$ star, its metallicity will resemble the sun with an abundance pattern as shown.

produced solely during supernova explosions. While Clayton's conclusion is somewhat sensitive to both an assumed model for Galactic chemical evolution and the unknown distribution of ^{26}Al interior to the solar orbit, we accept it for now and merely point out that revisions of reaction rates and stellar model have lead to a substantial increase in the supernova production of ^{26}Al over that estimated by Woosley and Weaver (1980), $^{26}\text{Al}/^{27}\text{Al} \approx 10^{-3}$. Here, for the new $25 M_{\odot}$ model, $^{26}\text{Al}/^{27}\text{Al} \approx 6 \times 10^{-3}$ and $6.9 \times 10^{-5} M_{\odot}$ of ^{26}Al are ejected. Of this ^{26}Al , a lesser fraction, $1.9 \times 10^{-5} M_{\odot}$ comes from proton capture on magnesium in the hydrogenic envelope (see also Dearborn and Blake 1985). Similarly, in the $15 M_{\odot}$ model, $7.1 \times 10^{-6} M_{\odot}$ was ejected in the hydrogenic envelope. These values would have been larger if the initial metallicity had been greater or if mass loss had been included in the calculation. On the other hand, this is an upper bound to the ^{26}Al that could be produced in 15 or $25 M_{\odot}$ stars of constant mass and solar metallicity since we considered all reactions on ^{25}Mg as going to the long lived ground state of ^{26}Al . In fact, some would go to the isomeric state (Harris et al. 1983) and be lost. Also recent re-examination of the $^{25}\text{Mg}(p,\gamma)^{26}\text{Al}_g$ reaction rate by Peter Parker at Yale (private communication) substantially reduces the value over that reported by Champagne, Howard, and Parker (1983) and included in the tables of Caughlan et al. (1985) for temperatures less than 40 million degrees (and increases it for higher temperatures). Typically hydrogen burns at temperatures at the cooler end of this range in 15 and $25 M_{\odot}$ stars and the synthesis of ^{26}Al in such stars will be correspondingly reduced. On the other hand, losing the ^{26}Al in a wind can increase its effective production in the hydrogen envelope since getting the material off of the star reveals the radioactivity to possible observation before a portion has had time to decay while waiting for the supernova (Dearborn and Blake 1985). More calculations are needed, especially using the modified $^{25}\text{Mg}(p,\gamma)^{26}\text{Al}_g$ reaction rate. Calculations are also needed of nucleosynthesis in nova ejecta enriched in neon and magnesium (see Truran, this volume). It is interesting that supernovae like the $25 M_{\odot}$ star described here, though perhaps still unable to account for all of the signal observed by Mahoney et al. should still give a substantial flux. We crudely estimate this to be about 10% of the HEAO-C signal (ignoring production in the hydrogen envelope). Since the Gamma-Ray Observatory will have at least 10 times the sensitivity of HEAO-C to these γ -lines (Kurfess et al. 1983), the signal from supernovae may still be discernable (depending upon the poorly determined Galactic distribution of the HEAO-C signal).

4. Type I Supernova Models - General Comments

There are many reasons for believing that Type I supernovae might be the result of accreting white dwarfs provoked into thermonuclear explosion. Observational evidence favors the association of such supernovae with a low mass population. They are not preferentially situated in the spiral arms of spiral galaxies (cf. Maza and Van den Bergh 1976) and they do occur in elliptical galaxies where no Type II supernovae are seen (Tammann 1977) and no young stars are expected (although see Oemler and Tinsley 1979). Type I supernovae, by definition, lack hydrogenic lines in their spectra as would be the case if a white dwarf exploded. The velocities inferred from spectral measurements of Type Is and the energies of the explosions agree with what one would obtain by converting a fraction of a white dwarf mass to iron (or ^{56}Ni). Further observational evidence supporting this inference is provided by the fact that iron is seen in the explosions (Kirshner and Oke 1975; Wu et al. 1983) as well as the radioactive decay product of ^{56}Ni , ^{56}Co (Axelrod 1980a,b; Branch 1984a,b). Furthermore the degenerate nature of a white dwarf guarantees that a nuclear runaway will convert a substantial fraction of its mass to iron on a short time scale (Hoyle and Fowler 1960) with the resulting light curve generated by the decay of these same radioactive species (Pankey 1962; Truran, Arnett, and Cameron 1967; Colgate and McKee 1969; Arnett 1979; Chevalier 1981; Weaver, Axelrod, and Woosley 1980). Finally, Type I supernovae are a very uniform class of events which might be understood if they all had a very similar origin, a compact object that creates 0.3 to 1.1 $M_{\odot}$ of ^{56}Ni (Axelrod 1980a,b). For all these reasons it is generally presumed that Type I supernovae are the consequence of the thermonuclear disruption of an accreting white dwarf star.

Generally speaking, there are two mechanisms for the explosion: *detonations* and *deflagrations*. Detonations have been recently reviewed by Nomoto (1982ab) and by Woosley, Taam, and Weaver (1985); deflagrations have been recently examined by Nomoto, Thielemann, and Yokoi (1984) and by Woosley, Axelrod, and Weaver (1984). See these reviews for extensive discussions of specific models and references to the published literature. Burning occurs by detonation when the runaway ignites in *helium-rich* material. This may occur either at the center of an accreting helium white dwarf or, off-center, at the base of an accreted helium layer on a carbon-oxygen or an oxygen-neon white dwarf. Ignition in helium occurs at sufficiently low density that burning to nuclear statistical equilibrium at $T \sim 8 \times 10^9$ generates a large local increase in the pressure, a factor of 5 being broadly representative. A higher temperature cannot be attained by the combustion to iron since photodisintegration truncates the energy deposition. This large overpressure leads to a shock wave. As material passes through the shock wave, typically at speeds of 2 to 3 times the speed of sound, the temperature and density rise, virtually instantaneously, and nuclear burning gives enough energy to keep the shock wave going. Because the expansion is supersonic,

material does not have time to "get out of the way" ahead of the burning front and essentially the entire layer of helium is converted to iron group elements, especially ^{56}Ni . Indeed, if the helium layer is thick enough and the density in the carbon-oxygen core not too high, the shock wave driven into the carbon oxygen core may propagate as a successful detonation wave as well in which case the star is totally disrupted and there is no bound remnant. More typically, a portion of the carbon oxygen core stays behind as a white dwarf remnant of the supernova.

The detailed computer models of this variety (referenced above) have been quite successful in reproducing the light curves observed in Type I explosions and in synthesizing, in proper solar proportions, some of the major isotopes of the iron group elements. Unfortunately this class of model is not consistent with spectroscopic restrictions. Ignition at a lower density and conversion of all the ejected material into iron gives an explosion that is both too energetic, i.e., the velocities are too high compared to observed line widths and Doppler shifts, and lacking in substantial quantities of intermediate mass elements, especially silicon and calcium, that are quite prominent in the early time spectrum. For this reason, greater attention is given, at least presently, to deflagrating models.

Deflagrations generally occur when an accreting carbon-oxygen white dwarf approaches the Chandrasekhar mass and ignites in its center. Typical ignition conditions require a balance of neutrino losses by the plasma process to balance the nuclear energy generation by a highly screened carbon fusion reaction and imply $\rho \sim 2 \times 10^9 \text{ g cm}^{-3}$. At this high density a temperature of $8 \times 10^9 \text{ K}$ gives only a small increment in the large Fermi pressure, $\leq 10\%$, which is not enough to generate a self-consistent detonation wave. Instead, once the burning front begins to travel, unburned fuel expands ahead of the front, having been notified by a sonic precursor of the events transpiring in the core. Density decreases as material crosses the burning front, generates heat, and expands. Perhaps most importantly, the expansion of a portion of the white dwarf is rapid enough that the (subsonic) burning front never overtakes it and thus unburned fuel is ejected, as well as a portion of the star that experiences intermediate burning temperatures and produces intermediate mass elements.

5. Problems with Deflagrations and Some Possible Solutions

The problems with deflagrating models center upon the poorly defined nature of the burning front and gross uncertainty in the velocity with which it propagates. The only clear limits are that it be subsonic yet lead, ultimately to the rapid expansion of the entire star. If the deflagration moves slowly compared to sound, then the expansion of the star is well underway before the burning front passes a fiducial point, say half-way out in mass. Thus less matter is burned to iron and, for starting points of similar gravitational binding, the supernova energy and velocity are smaller. Less

^{56}Ni produced also means a dimmer, briefer light curve, eventually, for $M_{\text{ss}} \leq 0.3 M_{\odot}$, too dim and too brief to be in accord with observations. A rapid deflagration velocity, on the other hand, gives the converse of these properties.

A generic problem with deflagrations is that, so far, they give unacceptable nucleosynthesis for the isotopes of the iron group. In the work of Nomoto, Thielemann, and Yokoi (1984), for example, $^{54}\text{Fe}/^{56}\text{Fe}$ is overproduced by a factor of 3.9 (Model W7) compared to the sun. Woosley, Axelrod, and Weaver (1984) find an even larger overproduction in a similar model. In this volume, F. Thielemann presents a recalculation of Model W7 that properly includes the " α -rich freeze-out" of Woosley, Arnett, and Clayton (1973). In the new study, he finds much of the material that previously was in the form of ^{54}Fe is transformed into ^{58}Ni . This helps things a bit since elemental ratios in the sun are more poorly known than isotopic ones, but the abundance of ^{58}Ni in the sun is comparable to that of ^{54}Fe and the overproduction is still quite large, about a factor of 5 in the new calculation for any reasonable value of electron capture rates.

The difficulty stems from the large ignition density of the deflagrating models which leads to a great deal of electron capture during the explosion. One would like to have almost no capture since adequate ^{54}Fe (or ^{58}Ni) can be created by just those neutrons available from conversion of the initial metallicity of the star to ^{22}Ne . In fact, $^{54}\text{Fe}/^{56}\text{Fe}$ resulting from nuclear statistical equilibrium with no electron capture is about 0.05 ($Z/0.015$), i.e., about solar for solar metallicity. Perhaps those white dwarfs that make Type I supernovae are metal deficient, but obviously the amount of electron capture must be kept to a minimum.

Two resolutions have been suggested for this dilemma (the second being reported here in published form for the first time). The first attempts to decrease the central ignition density by raising the accretion rate to a very high value. Iben and Tutukov (1984) and Webbink (1984) have discussed the possibility that two carbon oxygen white dwarfs might merge as a consequence of gravitational radiation in a binary pair. Such probabilities as one can legitimately assign to such uncertain events are consistent with observed SNI rates. In the final merger, the effective accretion rate must be very large; for lack of any calculation, an Eddington value is usually assumed. Unfortunately studies (thus far unpublished) by Nomoto and by us show that such high accretion rates do not, in the general case, lead to central carbon ignition. Compression of the outer layers by the rapid mass addition leads to off-center burning that gradually, on a thermal diffusion time, propagates inwards until the entire white dwarf has been converted to a oxygen-magnesium-neon composition. In the specific calculation leading to Figure 8, a carbon-oxygen white dwarf of $0.7 M_{\odot}$ and central temperature 5×10^7 K was allowed to accrete carbon and oxygen at a rate $2 \times 10^{-5} M_{\odot} \text{ y}^{-1}$. Once the mass had increased to $1.06 M_{\odot}$ carbon burning ignited in a shell at $0.77 M_{\odot}$. Over the next 19,000 years the carbon shell burned into the star at the same

Fig. 8 - Composition of a white dwarf that has experienced a very high accretion rate and off-center carbon ignition. The composition of an initially $0.7 M_{\odot}$ white dwarf of carbon and oxygen is sampled 37,000 years after the onset of accretion at $2 \times 10^{-5} M_{\odot} \text{ yr}^{-1}$. Very small regions of unburned carbon and oxygen still exist in the inner $0.02 M_{\odot}$ and outer $0.03 M_{\odot}$ of the star. The central density at this point is $2.57 \times 10^8 \text{ g cm}^{-3}$ as the white dwarf is nearing the Chandrasekhar mass. Continued evolution will lead to a neutron star.

time that the star's mass continued to increase eventually reaching the center about the time the Chandrasekhar mass was achieved. No explosive burning or dynamic events were observed. This star presumably does not ignite neon or oxygen burning until such high central densities are reached that electron capture causes a collapse to a neutron star (Miyaji et al. 1980). Thus while such objects are very interesting candidates for producing certain x-ray variables, they are hardly Type I supernovae. No radioactive ^{56}Ni will be produced and very little matter will be expelled.

More recently we have examined this same sort of model for larger initial white dwarf masses and larger accretion rates. So long as the time scale for reaching the Chandrasekhar mass, $(1.44 - M_i)/\dot{M}$, is less than the thermal diffusion time to the center of the star, the burning front will not have time to reach the center before a carbon deflagration ignites. Also, larger white dwarf masses diminish the effects of off-center compression relative to central heating. For a white dwarf of $1.37 M_\odot$ accreting at $10^{-5} M_\odot \text{ yr}^{-1}$, for example, we find that carbon ignites in the center. Accretion onto a $1.0 M_\odot$ white dwarf is under study. However, since these larger masses and accretion rates imply a central ignition density comparable to what one generally uses with a Chandrasekhar mass white dwarf anyway, part of the motivation for going to the high accretion rate is lost (although one does retain the possibility of making a Type I supernova in the fashion suggested by Iben and Tutukov). We prefer to solve the problem of abnormal nucleosynthesis in deflagrating models a second way, by a more careful consideration of the properties of the propagating burning front.

The velocity of a conductive deflagration is $v_b \sim \delta/\tau \sim (l_{mfp}c/\tau)^{1/2}$ where v_b is the velocity of the deflagration (flame) front, δ is its thickness, l_{mfp} is the mean free path of the electrons ($= 1/\kappa\rho$), τ is a typical burning time in the front, and c is the speed of the electrons which here are relativistic. For a density $\rho \sim 2 \times 10^9 \text{ g cm}^{-3}$, the opacity is quite low owing to the long path length of the degenerate electrons. At $T_9 = 1$, $\kappa = 2 \times 10^{-5} \text{ cm}^2 \text{ g}^{-1}$, and at $T_9 = 8$, $\kappa = 8 \times 10^{-3} \text{ cm}^2 \text{ g}^{-1}$. The nuclear burning time is of course very sensitive to the temperature which rapidly changes in the burning front. Typically half of the carbon burns, releasing about $2 \times 10^{17} \text{ erg g}^{-1}$, by the time the temperature has risen to 4 billion degrees. At this temperature, $\tau \sim 10^{-10} \text{ s}$ and thus, with obvious great uncertainty, $\delta \sim 10^{-3} \text{ cm}$ and $v_b \sim 10^7 \text{ cm s}^{-1}$. These values are interesting, but largely irrelevant since the flame will not propagate as a stable conductive front (cf. Müller and Arnett 1982).

For any reasonable conditions in the white dwarf star, the kinematic viscosity (Spitzer 1962; Landau and Lifshitz 1959) will be sufficiently small and the characteristic velocities and dimensions sufficiently large that the Reynolds number is very great. Thus the flow will be turbulent. Specifically, for $R > 10^5$, as noted by Zeldovich and Rozlovski (1947 as referenced on p. 477 of Landau and Lifshitz 1959), a spherically propagating combustion front in free space will have self-turbulence with characteristic dimension equal to the radius of the spherical flame. Thus we envision the runaway

initiating as a point in the center of the star, stabilized perhaps for some brief period by convection, but eventually localized to a small region. The pressure scale height as the runaway commences is 5.3×10^7 cm and contains $0.35 M_{\odot}$, but carbon burns at a rate less than the sound crossing time for this region once the temperature exceeds 1.5×10^9 K. As the temperature in the center continues to rise, ultimately reaching 8×10^9 K, the runaway becomes even more localized. Thus the flame front initially propagates in a very small region and, as we have just discussed, this means that the initial scale of the turbulence will also be quite small and the velocity of the flame relatively slow. As the burning front encompasses more and more material, the turbulent scale grows and so does the velocity. Thus one may approximate $v_b \sim kr_b$ where k is a constant of proportionality and r_b is the radius of the burning front. The solution to this equation is $r_b = r_0 \exp(kt)$ and hence $v_b = v_0 \exp(kt)$. Of course, one must impose the restriction that the turbulent motion remain subsonic, so as a maximum upper bound (which may, in fact, be approached), $v_b^{max} \leq c_s$, the sound speed. Thus we adopt

$$v_b(r) = fc_s (1 - e^{r/R_0})$$

where f and R_0 are adjustable constants, $f < 1$, R_0 less than the radius of the (expanding) white dwarf, which, in practice, will be determined by requirements on explosion energy, light curve, and nucleosynthesis.

This prescription is obviously arbitrary, but we feel that it relates more physically to the actual circumstances than any formulation based upon a convective mixing length, i.e., pressure scale height as used in all previous studies (cf. Nomoto, Thielemann, and Yokoi 1984; Woosley, Axelrod, and Weaver 1984). Certainly in its initial stages, the burning front has little knowledge of the pressure scale height. It propagates in a region of nearly constant density and pressure. The relevant scale is the scale of the turbulence.

We are quite concerned about the very early propagation of the front because the isotopic nucleosynthesis is so sensitive to it. A white dwarf near the Chandrasekhar mass obviously has $\Gamma \approx 4/3$ and is only marginally stable against large excursions in radius either inwards or outwards. Thus burning only a tiny fraction (we estimate $\leq 0.03 M_{\odot}$) of carbon and oxygen in the center of the star releases enough energy that it would expand to the point where its central density has dropped by a factor of 10. What is needed then, in order to avoid the excess electron capture that has plagued the deflagration model, is only that the star have time to respond to this burning and expand before the burning front encompasses a substantial fraction of its mass. Later, when the expansion has occurred, one then wants the burning front to go very fast in order to burn enough of the star that an energetic explosion results. This is the behavior that equation (1) will give and seems, to us at least, to be what is physically expected. We are carrying out a calculation to see if it will all work quantitatively.

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