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Gravitational collapse versus thermonuclear explosion of degenerate stellar cores

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Abstract

In this paper we review the behavior of growing stellar degenerate cores. It is shown that ONeMg white dwarfs and cold CO white dwarfs can collapse to form a neutron star. This collapse is completely silent since the total amount of radioactive elements that are expelled is very small and a burst of γ-rays is never produced. In the case of an explosion (always carbon-oxygen cores), the outcome fits quite well the observed properties of Type Ia supernovae. Nevertheless, the light curves and the velocities measured at maximum are very homogeneous and the diversity introduced by igniting at different densities is not enough to account for the most extreme cases observed. It is also shown that a promising way out of this problem could be the He-induced detonation of white dwarfs with different masses. Finally, we outline that the location of the border line which separetes explosion from collapse strongly depends on the input physics adopted.

Dans cet article on revise le comportement d'un noyau stellaire dégénéré qui grandit. On montre que les naines blanches d'ONeMg et celles de CO, froides et massives, peuvent s'effondrer pour former une étoile à neutrons. Cet effondrement est complètement silencieux puisque la quantité totale d'élements radioactifs expulsée est très petite et on ne produit pas d'eruption de rayons gamma. Dans le cas d'une explosion (toujours pour des noyaux de carbone-oxygène), le résultat des calculs reproduit assez bien les propriétés observées des supernovae de Type Ia. La courbe de lumière et les vitesses correspondant au maximum sont très homogènes et les variations introduites par des différences dans les densités d'ignition, ne suffisent d'expliquer les cas observés les plus extrèmes. On montre aussi qu'une solution prometeuse à ce problème pourrait être la détonation induite par l'ignition de l'helium á la surface de naines blanches de masses assez differentes. Finalement, on signale que la frontière entre l'explosion et l'effondrement depends fortement de la physique que l'on introduit.

9.1 Introduction

Exploding stars play a very important role in the evolution of galaxies since they inject about 10^{51} ergs of mechanical energy per event to the interstellar medium. They are at the origin of the neutron stars and, consequently, they are responsible of the existence of pulsars and galactic X-ray sources. Finally, during the explosion they inject several solar masses of newly synthesized elements, which completely shapes the chemical evolution of the galaxies. In all cases, a stellar core, supported by the pressure of degenerate electrons is involved. The reason is twofold:

a) At high densities, the electron pressure is dominant. This pressure is composed by two terms, a leading one, only depends on the density alone, and another one that vanishes whith temperature, i.e.:

$$P_e = P_0(\rho) + P_1(\rho, T) \tag{9.1}$$

$$P_1 \to 0\, T \to 0$$

Under degenerate conditions P_0 is dominant. If thermonuclear reactions start at some point, matter cannot expand in order to control them and a thermonuclear runaway that incinerates all the matter to iron peak elements occurs. As a consequence of the high temperatures resulting from the process, a burning front appears and propagates through all the star.

b) If electrons are not relativistic, the leading term is $P_0 \propto \rho^{5/3}$, and from the dimensional point of view it scales like $P \sim M^{5/3}R^{-5}$. If electrons are relativistic, $P_0 \propto \rho^{4/3}$ and it scales as $P \sim M^{4/3}R^{-4}$. The condition of hydrostatic equilibrium requires a functional dependence $P \sim M^2R^{-4}$. Thus, in the extremely relativistic case there is not a definite scale length and, depending on how fast the energy is injected or removed, it is possible to obtain an explosion (observed as a supernova) or a collapse (to form a neutron star). In the case of stellar cores, the outcome depends on the rate at which energy is injected by the burning front as compared with the rate at which it is removed by the electron captures on the burned

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material. The velocity of the burning front is a function of the density, temperature and chemical composition. The rate of electron captures on the incinerated material, which is essentially made of ^{56}Ni, is a function of the Fermi energy. Therefore, the outcome -explosion or collapse- depends on the density at which the fuel is ignited, for a given velocity flame.

9.2 Physics of the explosion

9.2.1 The thermonuclear runaway

The basic condition to obtain an explosion is that enough energy must be released before matter can dynamically react. In stars, this condition can only be fulfilled in electron degenerate structures. The reason is that if pressure were dominated by the classical ideal gas law, any energy release would produce an adiabatic expansion that would quench the reaction. In the degenerate case, the energy generation rate will increase because of the temperature increase and the process will be reinforced. Eventually, electrons will become nondegenerate and matter will start to expand. At this point, however, temperature will be so high that energy will be released in a time short as compared with the hydrodynamic time and a violent expansion will ensue.

The time necessary for the system to react to an overpressure is equal to the time spent by a sound wave to cross the system. Thus, the hydrodynamic time can be defined as:

$$\tau_{hd} \simeq \frac{h}{c_s} \tag{9.2}$$

where h is a characteristic dimension and c_s is the sound velocity. Since degenerate cores are in hydrostatic equilibrium, the balance ensures that the hydrodynamical time is of the order of the free fall time:

$$\tau_{ff} = \frac{1}{\sqrt{24\pi G < \rho >}} = \frac{444}{\sqrt{< \rho >}} \tag{9.3}$$

i.e. the natural time scale for the free fall of a uniform pressureless self-gravitating sphere. Thus, the condition for instability is $\tau_n \leq \tau_{hd}$, where the characteristic nuclear heating time is defined as:

$$\tau_n = \frac{c_v T}{\dot{\epsilon}} \tag{9.4}$$

c_v is the specific heat at constant volume and $\dot{\epsilon}$ is the rate of energy generation. The temperature at which $\tau_n = \tau_{hd}$ is called the deflagration temperature and it should not to be mistaken for the ignition temperature, defined

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as the temperature at which neutrino losses balance the energy produced by thermonuclear reactions.

Typical stars burn their fuel quiescently because they stabilize their temperature below the deflagration temperature by means of adiabatic expansion. The efficiency of adiabatic cooling can be defined as the relative density change necessary to restore the pressure equilibrium upon the release of some quantity of energy during an adiabatic expansion (Mazurek and Wheeler, 1980). If the gas is ideal, the adiabatic expansion efficiency is:

$$\frac{\delta\rho}{\rho_0} \propto \frac{Q}{kT} \tag{9.5}$$

where Q is the energy released per reaction. Since Q is of the order of MeV, the cooling is very efficient for temperatures below 10^9 K ($\simeq 0.1 MeV$). If the electron component is strongly degenerate

$$(\frac{\partial P_e}{\partial T})_\rho \ll (\frac{\partial P_i}{\partial T})_\rho \tag{9.6}$$

where P_e and P_i are the electronic and ionic pressures respectively and

$$\frac{\delta\rho}{\rho} \propto \frac{Q}{kT}\frac{P_i}{P_e} \tag{9.7}$$

and adiabatic cooling is only efficient in the region where $P_i \geq P_e$.

9.2.2 Propagation of the burning front

When the thermonuclear fuel is ignited, burning propagates through all the star driven by the following mechanisms:

a) *Spontaneous burning.* In a contracting core, the condition for burning can be reached quasi simultaneously in several points. Therefore, due to the existence of an initially small temperature gradient, the burning can spread over a wide region without any transport mechanism (Blinnikov and Khokhlov 1986; Woosley 1986). The time scale for acceleration of the nuclear burning is:

$$\tau_{nuc} = \frac{\dot{\epsilon}_{nuc}}{(d\dot{\epsilon}_{nuc}/dt)} \tag{9.8}$$

and the location of the burning front changes with a phase velocity that is given by $v_{ph} = (d\tau_{nuc}/dr)^{-1}$. This velocity increases when the absolute values of the temperature and density gradients decrease. Therefore, regions with $v_{ph} \geq c_s$ will ignite spontaneously and the burning front will propagate supersonically there.

b) *Detonation.* In order to describe the properties of a burning front

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it is assumed that the fuel and the burned material are separated by a region, where the reactions take place, of a width δ much smaller than any characteristic length, l, of the system. If $\delta << l$ it is possible to connect both sides of the front by means of conservation laws of mass, momentum and energy. In the reference frame associated to the front, these equations can be written as (Landau and Lifchitz 1959, Mazurek and Wheeler 1980):

$$\rho_1 u_1 = \rho_0 u_0 \tag{9.9}$$

$$P_1 + \rho_1 u_1^2 = P_0 + \rho_0 u_0^2 \tag{9.10}$$

$$\epsilon_1 + \frac{P_1}{\rho_1} + \frac{U_1^2}{2} = \epsilon_0 + \frac{P_0}{\rho_0} + \frac{u_0^2}{2} + q \tag{9.11}$$

that are similar to those describing shock waves except for the presence of the term q that represents the energy released by reactions. The subindexes 0, 1 apply to the unburned and burned material respectively and the remaining symbols have their usual meaning.

The mass flux crossing the front is defined as $j = \rho_0 u_0 = \rho_1 u_1$ and can be written as:

$$j^2 = -\frac{P_0 - P_1}{V_0 - V_1} \tag{9.12}$$

which implies that the mass flux across the front is determined by the ratio between the difference of pressures and specific volumes at both sides of the burning front. Real solutions must satisfy $(P_1 > P_0)(V_1 < V_0)$ or $(P_1 < P_0)(V_1 > V_0)$. The first solution corresponds to a detonation and the second one to a deflagration. From the equation of conservation of energy it is possible to write

$$\epsilon_0 + q - \epsilon_1 + \frac{1}{2}(P_0 + P_1)(V_0 - V_1) = 0 \tag{9.13}$$

which is called the detonation adiabat (the case q=0 is called the shock adiabat). This equation, together with that defining the mass flux determines the final state once the characteristics of the burning front have been specified. The physical meaning of the intersection between the adiabat and the mass flux is clear. A shock heats and compresses the material to a state (P_s, V_s) given by the intersection of the shock adiabat with the line defined by j^2. Because of the increase in temperature, material burns and reaches the state (P_1, V_1) defined by the detonation adiabat and the j^2-line intersection. Since $q \geq 0$, $P_1 < P_s$ and $V_1 > V_s$ which implies that a rarefaction is associated to the postchock burning.

The family of solutions obtained in this way, with j as a free parameter, has an extremum for which j and the front velocity are minima. This solution, called Chapman-Jouguet detonation, corresponds to the case where the j-line is tangent to the detonation adiabat. It has the important properties of being univoquely determined by the thermodynamic properties of the material (including q) and having a propagation velocity that is equal to the sound velocity of the burned material. All the remaining solutions, called strong detonations, move subsonically with respect to the burned material.

In stars, due to the spherical simmetry, material must be at rest at the centre. Therefore, the velocity has to decrease from some positive value behind the front to zero at the centre. This means that a rarefaction wave, moving at the sound velocity, must follow the detonation. Since strong detonations are subsonic respect the burned material, they are overtaken by the rarefaction wave and only the Chapman-Jouguet one can survive.

Although it has not been completely elucidated, it is generally accepted that due to the high densities existing in the central regions of CO white dwarfs, the overpressures induced by the burning front cannot give rise to a detonation (Mazurek, Meier and Wheeler 1977). However, as the subsonic flame propagates into regions of progressively decreasing density, it accelerates and it can eventually become a detonation (Blinnikov and Khokhlov 1986; Woosley 1986). It is interesting to notice here that a detonation does not necessarily produces the complete incineration of the material into iron. If the density is smaller than $\simeq 10^7$ g/cm^3, nuclear statistical equilibrium cannot be reached and, as a consequence, elements of masses intermediate between C-O and Fe are produced. If the density is smaller than 10^6 g/cm^3 even the fuel, the C-O mixture, has no time to be exhausted (Khokhlov 1989). In the case that the fuel is He this is not true and Fe is synthesized.

c) ***Deflagration.*** Deflagrations are the solution to the conservation laws that fulfill the condition $(P_1 < P_2)$ $(V_1 > V_2)$. The main dificulty with these solutions arises from the fact that matter is subsonic at both sides of the front and boundary conditions behind it can affect the front as well as matter ahead. In the case of a star with a subsonic burning front propagating outwards with a velocity D, the condition of matter being at rest at the center demands the existence of a shock precursor that boosts matter outwards (Mazurek and Wheeler 1980) and causes the expansion of the star.

There are two modes whereby a nuclear deflagration can propagate inside a degenerate core: The laminar mode and the turbulent mode.

In the laminar case (often known as conductive front), electrons transport the energy released in the burning regions to the surroundings, inducing their ignition. The velocity can be estimated in the following way (Landau and Lifchitz 1959): The width of the front is given by $\delta \simeq \sqrt{\chi\tau}$, where χ is the thermometric conductivity and τ is the time that the burning lasts. The velocity of the burning front is thus given by D $\simeq \delta\tau = \sqrt{\chi/\tau}$. This velocity is always of the order or smaller than $10^{-2}c_s$, where c_s is the sound velocity. Useful approximations are provided by (Timmes and Woosley 1992):

$$v_{cond} = 92.0(\frac{\rho}{2\,10^9})^{0.805}(\frac{x_{12}}{0.5})^{0.889} \tag{9.14}$$

valid in the range $0.01 \leq \rho_9 \leq 10$ in the case of a C-O mixture and:

$$v_{cond} = 51.8(\frac{\rho}{6\,10^9})^{1.06}(\frac{x_{16}}{0.6})^{0.688} \tag{9.15}$$

valid in the range $1 \leq \rho_9 \leq 14$ in the case of an ONeMg mixture.

In the turbulent case, a hot and less dense layer is formed below a cold dense layer. Because of gravity, the interface is Rayleigh-Taylor unstable and both layers are mixed, that noticeably increasing the propagation velocity due to the increase in the efficiency of the conductive transfer. Although at present there is not a satisfactory theory of turbulent flames in stellar interiors, it is possible to make an estimate of the velocity of a turbulent flame in the central regions of a star. Nomoto et al (1984) proposed from mixing length arguments, a turbulent velocity $v_t \simeq \sqrt{g_{eff}l/2}$ where $g_{eff} = GMr^2/\delta\rho/\rho$ is the effective acceleration, $\delta\rho$ is the difference of densities between both sides of the front and l is the mixing length taken as $l = \min(r, \alpha H_p)$ where α is an adjustable parameter of the order of unity. The use of the mixing length theory has been questioned not only because of the jump in density across the front but also because all other characteristicsof the burning strongly violate the basic hypothesis of the theory itself. Woosley(1990) proposed a fractal description of the burning front to take into account its wrinkling and Livne and Arnett (1993) proposed to treat the turbulent deflagration in terms of an ablative front in order to correctly handle the growth of the different unstable modes. At present, the question is completely open and in fact the mixing- length model, despite being physically incorrect, provides for $\alpha = 0.7$ the best agreement with observations. In practice, numerical treatments like those of Sutherland and Wheeler (1984) and Unno (1967) are used. It must be also taken into account that electron captures behind the burning front can completely inhibit the development of the Rayleigh-Taylor instability (Timmes and Woosley 1992).

9.3 The collapse of degenerate cores

The evolution of the primary star in a close binary system can give rise to either a helium white dwarf, a carbon-oxygen white dwarf, or an oxygen-neon-magnesium white dwarf, depending on the initial parameters of the system. However, not all compositions can be involved in the accretion induced collapse. Helium white dwarfs can immediately be discarded. Their mass growth would lead to explosive helium ignition at the center of the star for a central density $\rho_c \leq 4\,10^8$ g/cm^3, no matter how low is the temperature (Sugimoto and Nomoto 1980). After incineration, electron captures would be too slow, the overpressures would be large enough to start a detonation and the star would thus be completely disrupted (Woosley and Weaver 1986).

This leaves only carbon-oxygen and oxygen-neon-magnesium white dwarfs as possible candidates for a core collapse. These groups have to be considered separately as accretion induced collapse poses different problems for each of them. Both cooling in the stage between the formation of the white dwarf and onset of mass accretion and reheating by mass accretion are specially relevant in the CO case, whereas semiconvection associated with electron captures plays an important role in the ONeMg case.

As it was stated in the Introduction, the collapse/ explosive behavior alternative for a degenerate core depends on the density at which the burning front starts, for a fixed velocity of the flame. Since there is not yet a theory giving the velocity of the burning front, it is not possible to decide which is the value of the density beyond which the collapse ensues. It is possible, however, to establish two firm bounds and a guess as to this critical value. The upper limit is determined by the maximum velocity of the flame: the sound velocity. In this case, the minimum density necessary to guarantee the collapse is $\rho \approx 2\,10^{10}$ g/cm^3. The lower limit is determined by the minimum velocity of the flame: the conductive velocity. In this case, the minimum density necessary to get a collapse is $8.5\,10^9$ g/cm^3. In a previous calculation by Canal et al (1992), this limit was set to $9.5\,10^9$ g/cm^3 because the Coulomb corrections to the equation of state of the ions were not correctly included. The guess can be obtained by using everal recepies appeared in the literature and making some kind of average. For densities higher than $9\,10^9$ g/cm^3 some of them predict a collapse, and for $9.5\,10^9$ g/cm^3 all of them do.

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9.3.1 The CO case

Carbon-oxygen white dwarfs can form in binary systems either by Roche lobe overflow just before or just after ignition of He in initially close binaries or by Roche lobe overflow during the early or the thermally pulsing asymptotic giant branch phases in initially wide binaries. In the latter case, common envelope evolution should allow enough orbital angular momentum to be lost so that the wide binary evolves into a close binary. An important question is the upper mass limit for CO white dwarfs formed in this way. Observations of classical novae give average masses of 1.23 $M_\odot$, although this figure is not truly representative of the average since several selection effects favor the detection of massive white dwarfs, and models for the recurrent nova U Sco give a mass $M \simeq 1.38\ M_\odot$ (Starrfield et al 1989). However, the most massive white dwarfs found in these systems may well be ONeMg and not CO white dwarfs. Theoretically, CO cores of single stars should ignite C non explosively when $M_{core} \geq 1.1 - 1.2 M_\odot$ (Iben and Tutukov 1986).

The behavior of the central layers of a growing CO core is determined not only by the local balance among the nuclear energy released, the neutrino losses and the compressional work, but also by the properties of the outer layers. If such CO core is part of an isolated white dwarf star, it will cool down because of the photospheric losses. If this core is growing, as is the case of an accreting white dwarf in a binary system, it is heated not only by the H and He burning shells but also by the compression of the outer, partially degenerate layers.

An important ingredient ofthe problem of determining the density at which central ignition happens comes from the fact that nonperturbed white dwarfs cool down and eventually solidify. For instance, a CO white dwarf with $M \geq 1 M_\odot$ starts to solidify after 1 Gyr since its formation. That implies that the adiabatic coefficient is small $(\partial \ln T / \partial \ln \rho)_s \simeq 0.5$, that the heating by compression is gentle (Hernanz et al 1988) and that, in the absence of the influence of the outer layers, the ignition of the carbon-oxygen mixture is entirely controlled by pycnonuclear reactions (Canal and Isern 1979). In that case, the density at which the runaway starts is $\rho_c \simeq 10_{10}$ g/cm^3 instead of $4\,10^9$ g/cm^3 typical of the fluid phase.

The energy released by the compression of the outer layers is, by far, the most important heating mechanism. It can roughly be approximated by (Nomoto 1982):

$$L_g/L_\odot = 1.4\,10^{-3}(T/10^7K)(\dot{M}/10^{-10}M_\odot/yr) \tag{9.16}$$

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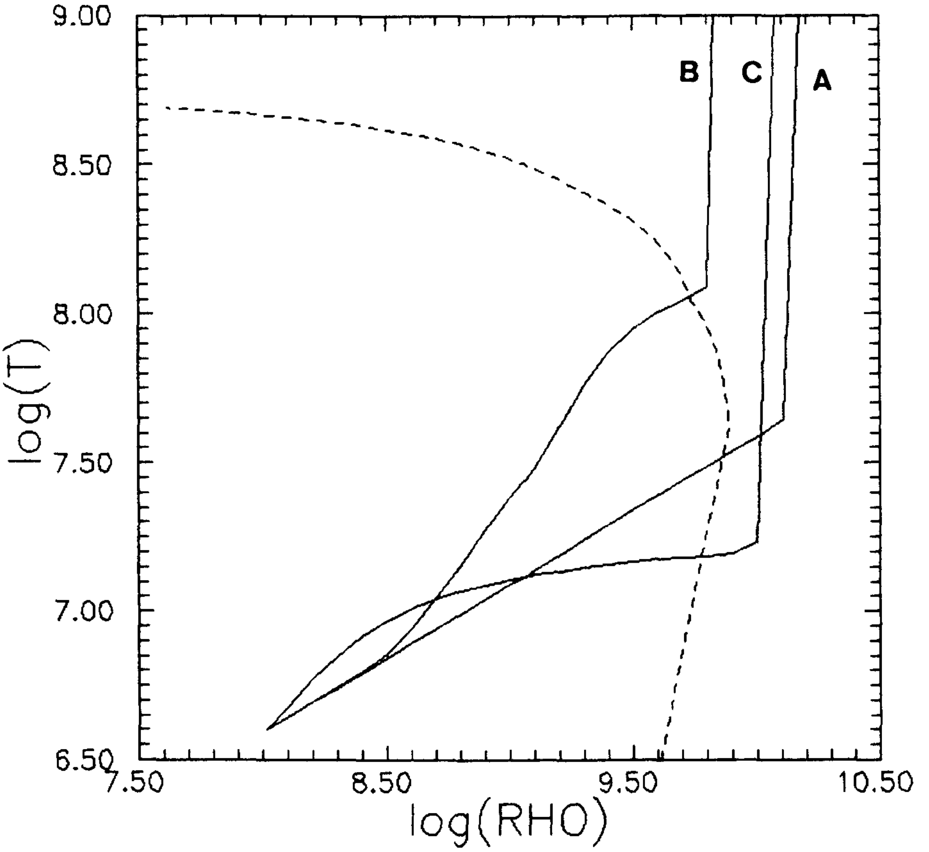

Fig. 9.1 Path followed by the center of a mass-accreting white dwarf with an initial mass of 1.15 $M_\odot$, $T_0 = 4\,10^6$ K and $x_c = x_0 = 0.5$. Cases A, B and C correspond tho $\dot{M} = 10^{-6}, 5\,10^{-8}$ and 10^{-10} $M_\odot$/yr respectively. The dashed line is the ignition line

If the accreted rate is smaller than $3\,10^{-10}$ $M_\odot$/yr, energy losses through the photosphere are dominant and the star cools down. If the accretion rate is higher, compressional heating is dominant and a thermal wave propagating inwards is generated. The consequence is that the inner layers are heated up, they cross the ignition line, defined by the condition that neutrino losses exactly balance the energy released by the carbon burning, in the (ρ, T) plane and a thermonuclear runaway starts. Nevertheless, if the initial mass of the white dwarf is $M \geq 1.2 M_\odot$ and the accretion rate is $\dot{M} \geq 5\,10^{-8} M_\odot/yr$, the thermal wave coming from the surface has no time to reach the center and the thermonuclear runaway is entirely determined by the local properties of matter (Hernanz et al 1988).

Figure 1 displays the evolution of the center of an accreting carbon-oxygen white dwarf. The three different behaviors already mentioned are clearly displayed: In case A, the thermal wave has no time to reach the central

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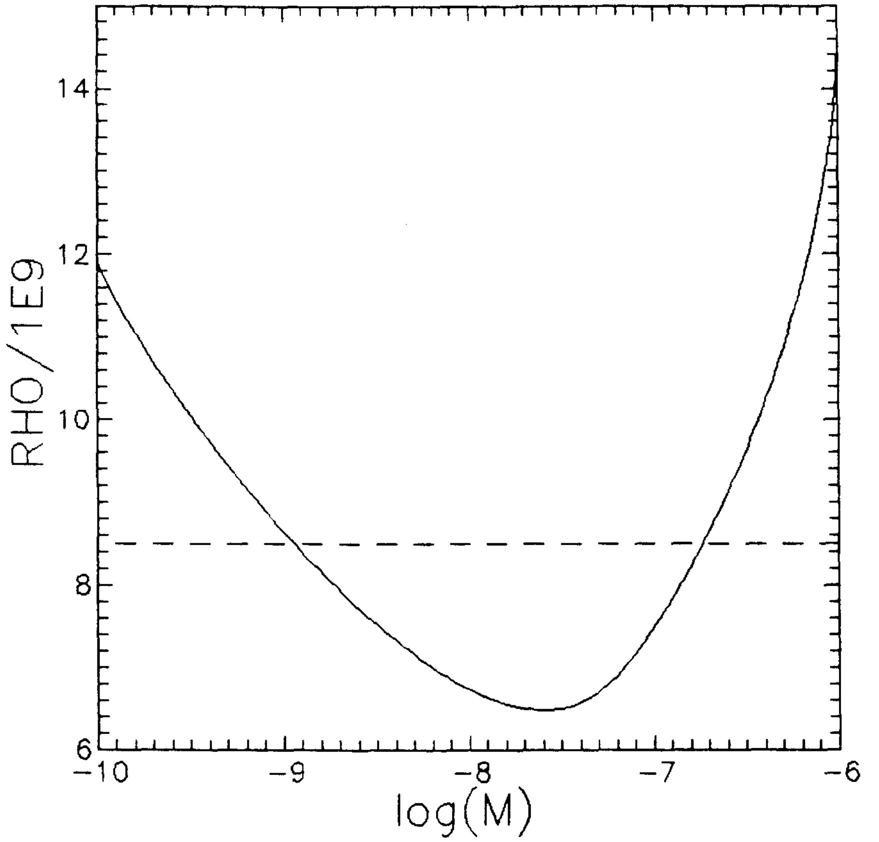

Fig. 9.2 Ignition density as a function of the mass accretion rate in the case of a white dwarf with the same characteristics as in Figure 1. The dashed line represents the minimum density necessary to have a gravitational collapse instead of a thermonuclear explosion

regions and the center follows a trajectory of slope 0.5 in the temperature-density plane. In case B, a strong thermal wave coming from the surface heats the material and finally induces its thermonuclear runaway. In case C, heating and surface cooling exactly balance each other after a transient phase and the white dwarf evolves isothermally.

Figure 2 displays the ignition density as a function of the mass accretion rate for a 1.15 $M_{\odot}$ white dwarf. Only those accreting either at a very small rate or at a high rate have a chance to collapse. Therefore, if we take into account that some of them can be initially hotter, it can be concluded that the majority of those white white dwarfs will explode and only a small fration will collapse. Notice also that if more massive white dwarfs were considered, the ignition density would increase and the collapse would be favored.

9.3.2 The ONeMg case

White dwarfs made of ONeMg would result, in close binary systems from loss of the helium layer when it expands to read giant size during C shell burning (Iben and Tutukov 1984; Nomoto 1984) and they are not expected to form from single-star evolution (Habets 1985). In contrast, their presence in close binary systems might be indicated by the observation of Ne-novae (Starrfield 1990; Truran and Livio 1986). They should be, on average, more massive and less frequent, by a factor $\simeq 10^4$, than CO white dwarfs (Iben and Tutukov 1984). When the ONeMg white dwarfs are compressed, the Fermi energy increases and nuclei undergo electron captures. The behavior of the temperature depends on the relationship:

$$\frac{dT}{dt} \propto (E_F + E_\gamma - E_\nu - E_{th,0}) \tag{9.17}$$

where E_F is the Fermi energy of electrons, E_γ is the energy of the excited states produced during the capture process, E_ν is the energy of the neutrino emiited and $E_{th,0}$ is the threshold energy for the electron capture to the ground state (Miyaji et al 1980). If the Fermi energy is high enough, the derivative is positive and the temperature increases in such a way that electron captures can trigger by themselves the thermonuclear runaway. The first element starting to capture electrons is ^{24}Mg ($\rho \simeq 4\,10^9$ g/cm^3), followed by ^{20}Ne ($\rho \simeq 9.1\,10^9$ g/cm^3) and ^{16}O ($\rho \simeq 1.9\,10^9$ g/cm^3).

Since electron captures produce an excess of entropy that tends to induce convection and, at the same time, a chemical gradient that tends to suppress it, the density at which the runaway actually starts is extremely sensitive to the way in which convection is handled (Mochkovitch 1984).

If the Schwarzschild criterion is applied, which does not take into account the presence of chemical gradients ($|\nabla| > |\nabla_{ad}|$ where ∇ is the actual gradient and ∇_{ad} is the adiabatic gradient), the entropy generated by electron captures induces the formation of a convective zone that transports very efficiently the excess of entropy. As the captures proceed, the star gradually contracts until a density of $2\,10^{10}$ g/cm^3 is reached. When this happens, electron captures on O trigger the ignition of this element and matter is completely incinerated to ^{56}Ni. Electron captures on ^{56}Ni are so fast that, independently of the speed of the burning front, the white dwarf collapses to a neutron star.

The Ledoux criterion takes into account the existence of chemical inhomogeneities. According to this criterion, the condition for the onset of convection becomes $|\nabla| > |\nabla_L|$, where ∇_L is equal to the adiabatic gradient plus a stabilizing term that depends on the gradient of the chemical

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composition (see, for instance, Cox and Giuli 1968). Consequently, convection is inhibited and strong local heating is produced. ^{24}Mg is exhausted and the temperature drops due to thermal neutrino emission before the temperature for explosive ignition is reached, but the laster happens at the onset of electron captures on ^{20}Ne. In this case, the ignition takes place at $9.2\,10^9$ g/cm^3 when the Takahara et al (1989) electron capture rates are adopted and the influence of the chemical potential on the electron threshold is properly taken into account.

9.3.3 The nonexplosive collapse of white dwarfs

Once nuclear reactions start at the center, the burning propagates through all the star. The one-dimensional calculations made up to now assume that the flame propagates with a velocity determined by the fastest mode of burning: either spontaneous, conductive or turbulent. The detonation mode is not considered since the central density, $\rho_c \geq 8\,10^9$ g/cm^3, of the models considered here is always very high. Except for minor differences, the behavior of the CO and ONeMg cores is always the same: burning propagates outwards, the electron captures reduce the mean electron mole number and induce the contraction of the star, the Chandrasekhar limit falls below the actual mass of the star, and finally the star collapses homologously.

It should be stressed here that in the case of CO white dwarfs the ignition happens in the interior of a solid. The strength of this solid is enough (Hernanz et al 1988) to prevent or at least strongly delay the developpement of convection. However, the neutrinos emitted by the electron captures on the burned material deposit enough energy to melt the crystal (the latent heat per nucleon is $l \simeq kT_m$, where T_m is the melting temperature). The energy deposited by neutrinos in their interaction with relativistic electrons is given by (Gehrstein et al 1976; Chechetkin et al 1980):

$$\epsilon_\nu = 2\,10^{15}\frac{E_\nu^3}{E_F}(\frac{2}{45}+\frac{2}{63}x+\frac{1}{140}x^2)Y_e\frac{L_{51}}{r_7^2} \tag{9.18}$$

where the units are erg/g/s, L_{51} is the neutrino luminosity in units of 10^{51} erg/s and r_7 is the radius in units of 10^7 cm, $x = E_\nu/E_F$, E_ν being the energy of neutrinos and E_F the Fermi energy in MeV. To quantify this situation, Isern et al (1990) considered a burning front placed at 10^7 cm from the center and assumed that $L_{51} = 0.1$, $E_F = 10$ MeV and $E_\nu = 9$ MeV. They found that the crystal melted in less than 1 second, a quantity that is of the order of the time necessary to develop the Rayleigh-Taylor instability.

The ejection of matter due to neutrino deposition following the collapse can be treated as a neutrino-driven wind (Hernanz et al 1993). A general description of these winds can be found in Duncan, Shapiro and Wasserman (1986) and in its relativistic form in Paczynski (1990). The equations that define the wind in its stationary form are:

$$4\pi r^2 \rho_0 v Y_\infty = \dot{M} \tag{9.19}$$

$$HY\dot{M} + L_\infty + L_\nu = \dot{E} \tag{9.20}$$

$$\frac{dP}{dr} = -\rho_0 H \frac{dlnY}{dr} \tag{9.21}$$

$$L_\infty = L(1 + v^2/c^2)Y^2 \tag{9.22}$$

$$Y = \sqrt{(1 - r_g/r)}/\sqrt{(1 - v^2/c^2)} \tag{9.23}$$

with, $r_g = 2GM/c^2$, $H = c^2 + (P + U)/\rho_0$, where $\dot{E}_0$ and $\dot{M}_0$ are the respective rates at which energy and mass are injected at the base of the wind, H is the enthalpy, ρ_0 is the rest mass density, L_∞ is the photon luminosity measured by an observer at infinity, L is the photon luminosity in the comoving frame, L_ν is the neutrino luminosity and k is the opacity. The equation of state is that of a gas composed by radiation, nuclei and e^-e^+ pairs, and the energy deposited by neutrinos takes into account the captures by protons and neutrons, the scattering by electrons and the creation of e^-e^+. Figure 3 displays an example that is in remarkably agreement with the numerical models of Woosley and Baron (1992).

The recent observations of γ-ray bursts by BATSE at the Gamma Ray Observatory have shown that these events are distributed isotropically but not uniformly in radius (Fishman et al 1991, Meegan et al 1991). These observations have opened the possibility of a cosmological origin. In this case, they must be placed at $z \approx 1$ on average and emit $\sim 10^{51}$ ergs in ~ 15 s (Paczynski 1991). One of the scenarios that have been proposed is the accretion induced collapse of a white dwarf (Dar et al 1992).

We have solved equations (19-23) for a set of reasonable values of the temperature and radius of the neutrinosphere. In all cases we have obtained a heavy wind characterized by: $\dot{M} \simeq \dot{E}(GM/R_\nu)^{-1} \geq \dot{E}/c^2$, where M and R_ν are the mass of the compact object and the radius of the neutrinosphere respectively. Since γ-rays can only emerge if the condition $\dot{M} < 10^{-2}(\dot{E}/c^2)$ is fulfilled (Paczynski 1990), we must conclude that the collapse of white dwarfs cannot explain the existence of γ-ray bursts (Hernanz et al 1993).

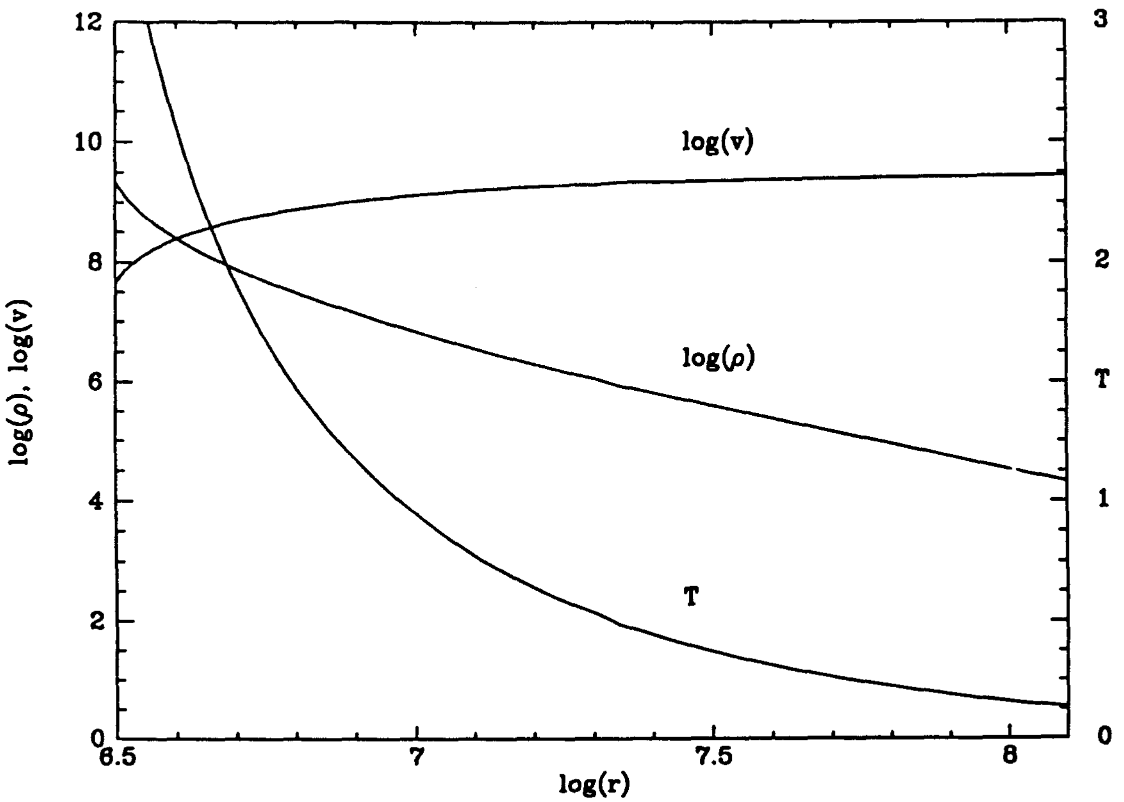

Fig. 9.3 Characteristics of a neutrino driven wind when $T_\nu = 4.96 10^{10}$ K and $R_\nu = 30$ km. At the sonic point $\rho_s = 9.95 10^5 \mathrm{g/cm^3}$, $T_s = 5.9 10^9$ K and $R_s = 210.4$ km. The intensity of the wind is $\dot{M} = 1.16 10^{31}$ g/s.

9.4 The explosion of degenerate cores

For a long time it has been assumed that the observational constraints that Type Ia supernovae should satisfy were the following ones:

1) The surfaces of their progenitors should be devoided of H at the time of explosion in order to explain the absence of Balmer lines in the spectra.

2) The progenitors should be long-lived stars in order to account for their occurrence in all types of galaxies, even the elliptical ones.

3) Their explosion should produce at least $\sim 0.5\ M_\odot$ of ^{56}Ni in order to account for the light curve and to explain the late time spectra.

4) Intermediate mass elements should be present in the outer layers in order to explain their spectra at maximum light.

5) The explosion should produce events with very homogeneous peak magnitudes (Miller and Branch 1990; Branch and Tamman 1992) whereas the light curve shapes and the photospheric expansion velocities might show some degree of variability.

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6) The abundance ratios of Fe and Ni isotopes should agree with the Solar System values after combining SNIa yields with those from gravitational collapse supernovae.

7) The death rates of the progenitors should agree with the observational estimates of the frequency of SNIa events.

Mainly due to points 1), 2) and 3) it is thought that SNIa are due to the explosion of CO white dwarfs in a binary system. Point 4) strongly suggests that the burning front propagates subsonically in the inner parts of the star and that only in the very outer layers, where $\rho \leq 10^7$ g/cm^3, it could propagate supersonically.

The observational situation mentioned in point 5) is very puzzling. It was claimed for some time that the rate of fading and the magnitude at the peak of the light curves as well as the expansion velocities of the photosphere of Type Ia supernova outbursts displayed a continuous behavior, i.e. the most luminous supernovae were declining more slowly and expanding more quickly than the less luminous ones (deVaucouleurs and Pence 1976; Pskovskii 1977; Branch 1981, 1982). This was challenged by Cadonau et al (1985), who examined the shape of twelve SNI light curves in a sample of elliptical galaxies and reached the conclusion that the dispersion of the light curves was smaller than 0.3^m when only photoelectric photometry was taken into account.

Concerning the maximum of the light curve, Miller and Branch (1990) examined the Pskovskii's sample and found that the dispersion in maximum brightness is smaller than 0.4^m if the inclinations of the galaxies are taken into account. Branch and Tammann (1992) proposed an absolute blue magnitude at maximum of $M_B = -19.6 \pm 0.4$. Concerning the postpeak decline, it has been shown that contamination by the light of the background galaxy and the K-corrections (Boisseau and Wheeler 1991; Leibundgut et al 1991) might account for the dispersion. Nevertheless, there are clear evidences of SNIa displaying peculiar behaviors. SN1885A, in M31, was very fast (deVaucouleurs and Corwin 1985). SN1986G in NGC128 (Phillips et al 1987) seems to have been intrinsically dim, to have a low expansion velocity and a fast light curve decline. SN 1991T (Filippenko et al 1992; Ruiz-Lapuente at al 1992) seems to have been overluminous and extremely peculiar in several aspects. Finally, the existence of different expansion velocities near the maximum of the light curve has been confirmed in a number of cases (Branch 1987; Branch et al 1988; Schneider et al 1988; Philips et al 1987; Barbon et al 1990). These velocities range from 10,000 km/s in SN1986G, SN1986A and SN1989B to at least 15,000 km/s in SN1983G and SN 1984I (Branch and Tammann 1992). The most extreme case of peculiar behavior has been provided by SN1991bg in NGC 4374 (an elliptical galaxy in the Virgo cluster) which was clearly underluminous: at maximum light its B magnitude was $\sim 2.5^m$ fainter and its V magnitude $\sim 1,6^m$ fainter than a normal SNIa in the same galaxy, it declined very fast after maximum and entered the nebular phase sooner than other SNIa (Filippenko et al 1992; Leibundgut et al 1993).

Therefore, the question to elucidate is the following one: is there a bulk of very homogeneous events, with some "dissidents" which can be explained just by allowing some minor changes in the main parameters of the deflagration/detonation model or is it necessary to build a new scenario to account for the existence of the "anomalous cases"?. The discovery of SN1991bg seems to point towards the second alternative.

Point 6) has been recently analyzed by Bravo et al (1993) and the combined yields of SNI and SNII seem to account fairly well for the observed abundances.

Point 7) has recently turned out to be critical. A popular scenario for SNIa explosions involves the merging of two CO white dwarfs in a binary system due to the emission of gravitational radiation (Iben and Tutukov 1984). However, the negative results of searches for progenitor systems have raised serious objections to this scenario (Munari and Renzini 1992). This has renewed the interest in the single degenerate scenario (Wheelan and Iben 1973) where a white dwarf grows to the point of explosive ignition by accreting matter from a nondegenerate companion, typically a red giant or a supergiant, probably forming a symbiotic star. New estimates of the space density of symbiotic stars have increased their inferred numbers by a factor $\sim 10 - 100$. That means that the fraction of such stars that should reach the Chandrasekhar mass is just $\sim 4\%$ (Munari and Renzini 1992) or $\sim 40\%$ (Kenyon et al 1993) according to the higher and lower estimates respectively.

The observation of transient hydrogen lines in at least two SNIa, provides aditional support to the symbiotic scenario. In this context, it is especially relevant the discovery of H-lines in the nebular spectrum of SN1991bg (Ruiz-Lapuente et al 1993), which can be interpreted as due to H-rich material, stripped from the companion by the kinematic interaction with the supernova ejecta, that appears as low velocity material in the late type spectra (Chugai 1986). Nevertheless, it is necessary to keep in mind the peculiar behavior of this supernova before generalizing this observational evidence.

The symbiotic scenario also opens up a new and interesting possibility. After burning, the accreted hydrogen is converted into heliun and accumulates on the surface of the star. Depending on the accretion rate and on the initial mass of the star, helium eventualy detonates and produces an inwards shock wave whose strenght increases due to geometric effects. Two dimensional hydrodynamic simulations (Livne and Glasner 1991) shows that this shock wave ignites a detonation in the center of the CO core that completely disrupts the star. This has lead the suggestion (Woosley and Weaver 1993, Canal 1993) that the helium detonation of CO white dwarfs with initial masses in the range 0.5 - 1.3 $M_\odot$ after accreting $\sim 0.1 - 0.2\ M_\odot$ of hydrogen-rich material could be at the origin of SNIa. Therefore, the question is whether both models, central ignition or He-induced detonation of CO white dwarfs, can coexist or one of them must be eliminated. Notice that the symbiotic scenario is compatible with the central ignition model and that one of the criteria to discriminate among them is their ability to reproduce the emergent variety of SNIa events.

9.4.1 Models igniting carbon at the center

This family of models assumes a mass accreting CO white dwarf in a binary system that approaches to the Chandrasekhar's mass and is partially incinerated by a subsonic burning front. A close examination of the preexplosion evolution reveals that the thermal runaway can happen at a density ρ in the range $2\,10^9 \leq \rho \leq 1.3\,10^{10}$ g/cm^3, that depending on the history of the binary system (Hernanz et al 1988). The shape of the light curve and the luminosity at maximum depend on the kinetic energy as well as on the total amount of ^{56}Ni newly synthesized and on its distribution accross the star, since γ-rays emitted by radioactive nuclei must be thermalized before escaping to contribute to the optical light curve. Both properties are affected by the total amount of electron captures undergone by the incinerated material, which depends on the ignition density, and by the total amount of matter that is completely incinerated (Graham 1987; Canal et al 1988).

For a given chemical composition, there are two properties that can modify the characteristics of the light curve. One is the velocity of the burning front and the other is the density at which the thermal runaway starts. If the burning front propagates as a Chapman-Jouguet detonation, the velocity of the burning front is completely determined by the thermodynamic properties of the burned material or, equivalently, by the density of the white dwarf. If the front propagates as a deflagration, the flame velocity is not uniquely determined by the density and it can be very different from one event to another, even if all the objects have a similar structure. In this section we examine the dependence on the density of the explosion characteristics.

The initial phases of the explosion, those encompassing from the ignition at the center to homologous expansion, have been modelled by several groups. The most critical, point as mentioned above, is the treatment of the burning front. A typical model is, for instance, that of Bravo et al (1993), who used an explicit difference scheme similar to that of Colgate and White (1966), an equation of state for the ion component taken from Ichimaru et al (1988) and for the electron component an ideal Fermi gas plus electron-positron pairs. Radiation was also included. The electron capture rates were taken from the compilation of Fuller et al (1982) and when not available there they were computed from the gross theory of $\beta-$decay (Kodama and Takahashi 1975). The equation of state for the nuclear statistical equilibrium (NSE) material has been computed using a set of 722 nuclei with $0.39 \leq Z/A \leq 0.50$. Mass excesses and partition functions were obtained from Thielemann (private communication). Nuclear burning was followed with an $\alpha-$network plus the 3α, C+C and O+O reactions. Si-burning was simulated by $^{28}\mathrm{Si}+7\alpha \rightarrow ^{56}\mathrm{Ni}$ and the rates were taken from Caughlan and Fowler (1988). NSE was assumed when $T > 2\,10^9$ K or $T > 5.5\,10^9$ K, for $\rho > 7\,10^7$ g/cm^3 and $\rho > 10^6$ g/cm^3, respectively.

It is also possible to obtain a rough bolometric light curve using the diffusion approximation and a finite difference scheme similar to that proposed by Falk and Arnett (1977). A flux limiter, of the form proposed by Alme and Wilson (1974), has to be used to avoid overestimate of the radiative flux in the outer layers. The equation of state is that of an ionized ideal gas plus radiation and the degree of ionization can be computed from the Saha equation (see Hoflich et al 1992 for details). The justification of the procedure only simplicity. Below 3500 K, the degree of ionization is very low. However, since the space is pervaded by energetic γ-photons, $E_\gamma \simeq 1$ MeV, coming from the radioactive nuclei, the degree of ionization of matter is higher than that corresponding to such temperature and the opacity is increased by several orders of magnitude. The resulting opacity is conveniently modelled by (Swartz 1991):

$$\kappa_c = \max[\kappa_{sc}, 1.4\,10^{-10}(\frac{\epsilon}{\rho})^{1/3}] \tag{9.24}$$

where κ_{sc} is the opacity obtained assuming local thermodynamic equilibrium, ϵ is the radioactive energy locally deposed and ρ is the density. The sources of opacity that have to be be considered are scattering Thomson by free electrons, bound-free and free-free transitions, and the contribution from the lines modified by the expansion effects (Karp et al 1977). The total average Rosseland opacity obtained in this way lies in the range of 0.05 to

0.1 cm^2/g. Finally, the energy deposited by γ-photons can be handled in several ways, the simplest one being to treat them as a simple absorption process (Sutherland and Wheeler 1984) assuming a γ opacity $\kappa_\gamma = 0.03$ cm^2/g, which is accurate enough for the majority of purposes.

The development of the Rayleigh-Taylor instability associated with the deflagration can be handled in several ways, no one being completely satisfactory. Table 1 displays the main characteristics of several models. The columns have the following meanings: ρ_9 is the density at which the central runaway starts, in units of 10^9 g/cm^3; M_b is the total burned mass, incinerated plus partially burned; M_{Ni} is the ejected mass of radioactive Ni; K_{51} is the kinetic energy in units of 10^{51} ergs, and M_{Fe} is the mass of ^{56}Fe synthesized. In the models labelled R, the development of the Rayleigh-Taylor instability associated with the deflagration was computed in the way proposed by Sutherland and Wheeler (1984), taking $\alpha = 0.7$. Models labelled J were designed to allow the burning front to propagate as a deflagration in the central regions and as a detonation in the outer layers. The development of the Rayleigh-Taylor instability was handled, in this case, with Unno's theory of time dependent convection (Unno, 1967). Concerning the two free parameters of the theory, the excess of temperature and the initial velocity of the "bubbles", ΔT_0 and v_0 respectively, it was assumed that $v_0 = 0$, and

$$\Delta T_0 = \frac{l}{2}[(\frac{\partial T}{\partial P})_S \frac{dP}{dr} - \frac{dT}{dr}] \tag{9.25}$$

where l is the mixing length. The last choices implies that the transition regime is very short and that the steady state is attained almost instantaneously. The characteristic mixing length scale was taken to be equal to the density length scale. In all cases, when the density was $\rho = 3 - 4\, 10^7$ g/cm^3 respectively, the deflagration spontaneously turned into a detonation that partially incinerated the material.

The models show that despite the increase in the total burned mass, which monotonically increases with the ignition density, the total mass of Ni produced during the explosion decreases and is strongly reduced in model R8. This is due to the electron captures on the incinerated material near the center, which are more important at high densities. Models J, however, propagate the burning faster, thus reducing the time available for electron captures, and so the final amount of ^{56}Ni is similar in all of them. The total kinetic energy dos not appreciably change with density because the increase of the burned mass is compensated by the energy losses due to electron captures and the initially greater binding energy. In the models J the total

Table 9.1. *General characteristics of the computed models*

Model	ρ_9	$M_b(M_\odot)$	$M_{Ni}(M_\odot)$	K_{51}	$M_{Fe}(M_\odot)$
R2	2.5	0.86	0.56	0.85	0.58
R4	4.0	0.89	0.52	0.91	0.59
R8	8.0	0.96	0.34	0.86	0.46
J2	2.5	1.00	0.63	1.42	0.68
J4	4.0	1.06	0.63	1.42	0.73
J8	8.0	1.19	0.51	1.40	0.66

Table 9.2. *General characteristics of the computed light curves*

Model	t_{bol}	M_{bol}	v_{ph}	β_0	β_1
R2	13	-19.23	10400	6.5	3.9
R4	13	-19.20	11000	6.5	4.0
R8	11	-18.81	11200	7.0	4.0
J2	12	-19.56	12424	7.1	3.3
J4	12	-19.54	12700	7.7	3.4
J8	10	-19.54	13800	8.5	3.4

amount of burned matter is very large and, consequently, the kinetic energy is also very large.

Table 2 displays the parameters that characterize the models. M_{bol} is the magnitude at maximum; t_{bol} is the time in days, from the explosion to the maximum; v_{phot} is the velocity of the photospheric layer, defined as the layer of optical depth unit, in km/s; β_0 and β_1 are the slopes in magnitudes per 100 days of the light curve 10 days after maximum and during the exponential tail.

The most noticeable fact in Table 2 is that the magnitude at maximum is the same in all the families of models and only the case R8 displays a significant deviation from the average. The reason is twofold: in the models where the runaway starts at high densities, the destruction of ^{56}Ni by electron captures in the central layers is counterbalanced in part by the production of this element in the outer layers and the proximity of the radioactive material to the surface, which facilitates the scape of photons. The large amount of ^{56}Ni newly synthesized and the small opacities of the models translate into very bright light curves, with the exception of the R8 one. Leigbungut et al (1991) have determined that the average apparent

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magnitude, free from extinction, of six SNIa in the Virgo cluster is 11.92 ± 0.11. If the distance to that cluster is assumed to be 20.03 ± 3 Mpc, their absolute magnitude would be $M_B = -19.6 \pm 0.4$ and only models J, i.e. the delayed detonation models, would be fully consistent with this value. However, if a distance of 16.5 Mpc is adopted (Jacoby et al 1992), the magnitude at maximum would be $M_B \simeq -19.2$ and models R would be completely acceptable.

Concerning the shape of the light curves and leaving aside the problem of their absolute callibration, models R perfectly fit the shape of the bolometric light curve since, on average, the observed rate of decline of the bolometric light curves is $\beta_0 = 6 \pm 0.5$ and $\beta \simeq 3.3$, although their exponential tail is too step. On the contrary, models J display a first decline after maximum and an exponential tail with a slower slope than in the case R. In any case, however, the light curves of each family are very similar and differ by less than 0.5^m, which implies that the ignition density does not seem to be able, just by itself, to introduce significant changes in the shape of the light curves.

The expansion velocity changes from model to model. In the family labelled R, the velocity differences are modest, less than 800 km/s. In the J models, this difference reaches 1500 km/s. In any case, these values define a range of variation that is smaller that the observed one. Due to the low opacities of the models, $\kappa \simeq 0.05 - 0.1$ cm^2/g on average, the photospheres of models J are rather deep and the velocity rather low, 11,000 to 12,000 km/s during maximum. These quantities are in agreement with the observations (Branch et l 1985), but there is a noticeable amount of matter, 0.01 $M_\odot$, moving at high velocities. Because of the lack of ultraviolet observations of supernovae, it is hard to obtain any conclusion from these figures. On one hand, Harkness and Wheeler (1990) obtained from UV observations that the maximum expansion velocity was 25,000 km/s for SN1981B. On the other hand, there are several supernovae like SN1984A (Wegner and McMahan 1987), SN1983G (McCall et al 1984), and SN1990N (Leibungudt et al 1991) that display material moving at very high velocities.

The nucleosynthesis can be computed with a post/processing code that uses the previously computed time evolution of the temperature and density for each shell. The most noticeable feature is the extraordinary increase of neutronized elements when the ignition density increases. This implies that the number of exploding high density white dwarfs which have contributed to the building of Solar System abundances, must had been very low in order to account for the observed values. It is also important to notice that even the models that ignite carbon at low densities produce an excess of

^{50}Ti, ^{54}Cr, ^{54,58}Fe, and ^{56}Ni. The first two nuclei are essentially produced in the central region and their overabundances cannot be avoided if, as it actually happens, the velocity of the burning front is very small at the center. Nevertheless, the degree of acceptability of these excesses can only be ascertained in the context of a model of the chemical evolution of the galaxy and taking into account the uncertainties that characterize our knowledge of isotopic and elementary abundances in stars and in the solar system. A simplified model (Bravo et al 1993) shows that the constraints introduced by the cosmic abundances of the elements are much less restrictive than previously thought.

9.4.2 Models igniting He off-center

The symbiotic scenario demands a non-violent burning of the freshly accreted hydrogen in order to effectively incorporate the newly synthesized helium onto the white dwarf. The usual limits for the behavior of hydrogen are as follow: Low rates, $\dot{M} \leq 10^{-9}$ M$_{\odot}$yr^{-1}, lead to a nova outburst that not only removes all the accreted matter but even erodes the original white dwarf. High accretion rates, $\dot{M} \geq 10^{-6}$ M$_{\odot}$yr^{-1}, lead to the formation of a red giant envelope, while intermediate rates in between these two limits lead to the formation of a common envelope.

The nova limit is based on calculations that assume spherically symmetric and soft accretion. Soft means here that material is deposited at the surface at rest and with the same entropy as the underlying material. It is thus cold and becomes strongly degenerate before igniting. Most likely, however, this material will form a disk around the compact star, and the accretion process will thus include angular momentum and kinetic energy dissipation. When those effects are taken into account, the actual range of $\dot{M}_H$ producing nova explosions becomes ill defined (Shaviv and Starrfield 1987, Sparks and Kutter 1987) and further research on this point is thus needed. Anyway, even in the most favorable cases, nova outbursts would limit the actual mass growth to at most 10% of $\dot{M}_H$.

The development of a common envelope (either due to accretion above the Eddington limit or to formation of a red-giant envelope should induce mass-loss by the system as a whole and thus inhibit further growth of the white dwarf. Nonetheless, accretion rates in the range $10^{-9} \leq \dot{M}_H \leq 10^{-6}$ M$_{\odot}$yr^{-1} allows to convert H into He trough steady combustion or weak flashes.

A further constraint is that the He layer resulting from the burning of the accreted H will explosively ignite only if it is accumulated at a rate $10^{-9} \leq$

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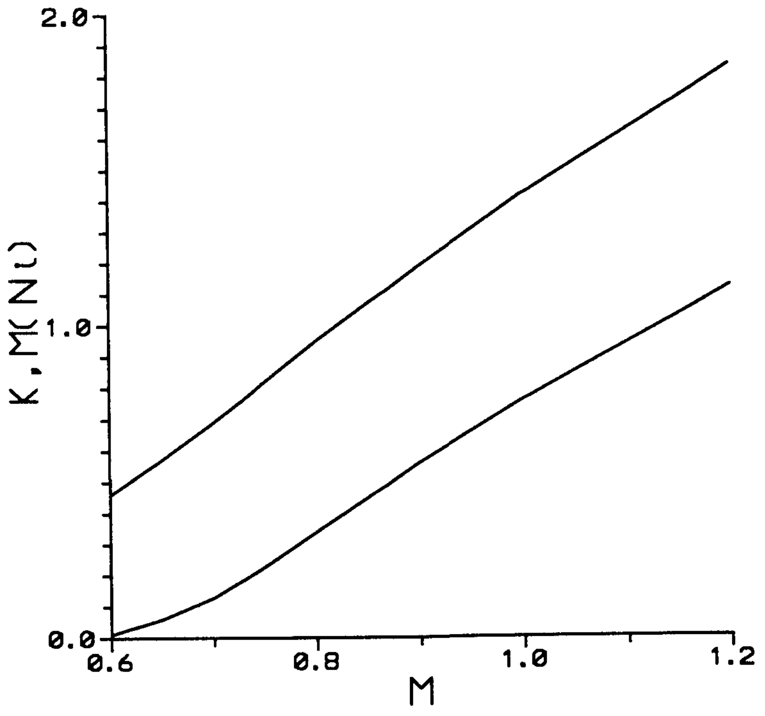

Fig. 9.4 Explosion energies, in units of 10^{51} ergs (upper line) and ^{56}Ni, in $M_\odot$ (lower line) of detonated white dwarfs

$\dot{M}_{He} \leq 5\,10^{-8}\ M_\odot yr^{-1}$. Above that limit it burns steadily (Nomoto 1982) and below this limit can accumulate safely depending on the parameters of the binary system. Of course, the same criterion applies if He is directly accreted from a companion (either degenerate or nondegenerate) that has previously lost its hydrogen envelope.

The ignition of He at the bottom of the envelope induces the formation of two strong shock waves. The first one becomes immediately a detonation that converts He into iron-peak elements, propagates outwards and sweeps out the outer envelope. The second one propagates inwards, increases its strength because of the spherical geometry and eventually turns into a detonation (Nomoto 1982). Figure 4 displays the kinetic energies and masses of ^{56}Ni synthesized by the detonation of low mass white dwarfs. The models, taken from Ruiz-Lapuente et al (1993), were computed with the implicit, one dimensional hydrocode described by Canal et al (1992). The light curves displayed in Figure 5 have been obtained assuming the diffusion approximation and a constant opacity $k = 0.2$ cm^2/g. It is clear that the range of

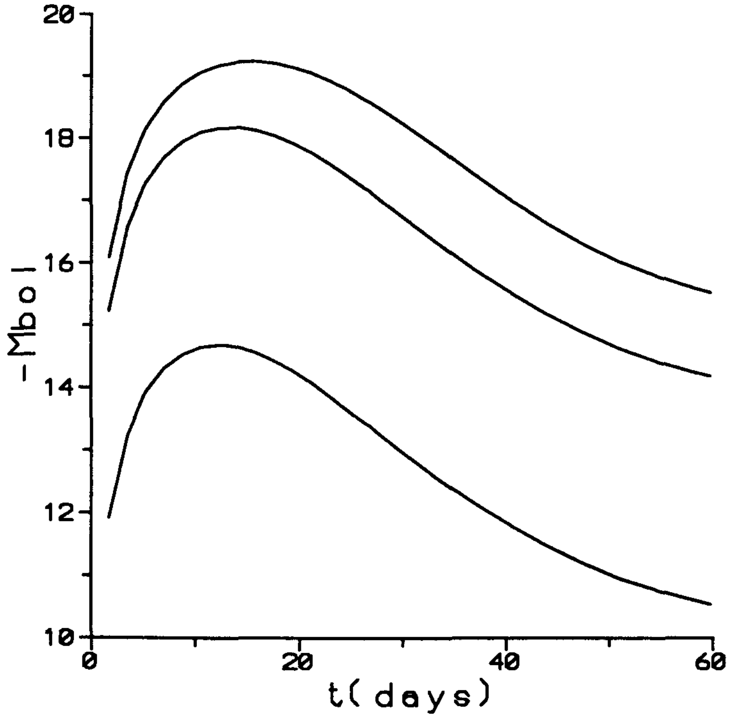

Fig. 9.5 Light curves obtained from the detonation of CO-white dwarfs of different masses. They correspond to the detonation of M^0_{WD} = 1.2, 0.8, 0.6 M⊙

variability that they show is higher than that obtained for central ignitions. Nevertheless, this problem still requires further study.

9.5 Conclusions

Mass accreting white dwarfs can either collapse or explode. The critical point is the density at which the thermonuclear runaway starts.

In the case of white dwarfs made of ONeMg, the ignition is triggered by electron captures on neon. The density at which it happens is $9.2\,10^9$ g/cm^3 and, despite the uncertainties on the properties of burning fronts, a collapse is almost guaranteed.

The ignition density of carbon-oxygen white dwarfs depends on the parameters of the binary system (initial masses and separation) which determine the instant at which mass transfer will start, its rate and chemical composition, as well as the initial mass of the white dwarf. If the white

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dwarf is massive enough ($M \geq 1.15\ M_{\odot}$) and cool enough, the thermonuclear runaway is delayed to densities higher than $9\,10^9$ g/cm^3 for both very high and very low accretion rates and a collapse ensues.

In both cases, the collapse is non explosive. The energy deposited by neutrinos in the outer layers only produces a heavy wind and there is not any γ-ray signal.

In the explosive case, the light curves and the photospheric velocities display some degree of variability due to the different ignition densities but that is not enough to account for the observations. One possible way out for this problem could be the induced detonation of CO white dwarfs with different initial masses.

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