

A call for a paradigm shift from neutrino-driven to jet-driven core-collapse supernova mechanisms

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ABSTRACT

Three-dimensional (3D) simulations in recent years have shown severe difficulties producing 10^{51} erg explosions of massive stars with neutrino-based mechanisms while on the other hand demonstrated the large potential of mechanical effects, such as winds and jets in driving explosions. In this paper, we study the typical time-scale and energy for accelerating gas by neutrinos in core-collapse supernovae (CCSNe) and find that under the most extremely favourable (and probably unrealistic) conditions, the energy of the ejected mass can reach at most 5×10^{50} erg. More typical conditions yield explosion energies an order of magnitude below the observed 10^{51} erg explosions. On the other hand, non-spherical effects with directional outflows hold promise to reach the desired explosion energy and beyond. Such directional outflows, which in some simulations are produced by numerical effects of 2D grids, can be attained by angular momentum and jet launching. Our results therefore call for a paradigm shift from neutrino-based explosions to jet-driven explosions for CCSNe.

Key words: stars: massive – supernovae: general.

1 INTRODUCTION

Eighty years after Baade & Zwicky (1934) first suggested that supernovae (SNe) are powered by stars collapsing into neutron stars (NS), the processes by which part of this gravitational energy is channelled to explosion remain controversial. Wilson (1985) and Bethe & Wilson (1985) refined the neutrino mechanism (Colgate & White 1966) into the delayed-neutrino mechanism, whereby neutrinos emitted within a period of ~ 1 s after the bounce of the collapsed core heat material in the gain region ($r \approx 100$ – 200 km). This subsequent neutrino heating was thought to revive the stalled shock thereby exploding the star and producing a canonical core-collapse supernova (CCSN) with an observed energy of $E_{\text{exp}} \gtrsim 1$ foe, where $1 \text{ foe} \equiv 10^{51}$ erg.

In the last three decades, sophisticated multidimensional simulations with increasing capabilities were used to study the delayed-neutrino mechanism (e.g. Bethe & Wilson 1985; Burrows & Lattimer 1985; Burrows, Hayes & Fryxell 1995; Fryer & Warren 2002; Buras et al. 2003; Ott et al. 2008; Marek & Janka 2009; Nordhaus et al. 2010b; Brandt et al. 2011; Hanke et al. 2012; Kuroda, Kotake & Takiwaki 2012; Mueller, Janka & Marek 2012; Bruenn et al. 2013, 2014; Mezzacappa et al. 2014; Müller & Janka 2014). The outcome of such numerical experiments varied widely with many failing to revive the stalled shock while others produced tepid explosions with energies less than 1 foe. Historically, in spheri-

cally symmetric calculations (1D), the vast majority of progenitors cannot even explode (Burrows et al. 1995; Rampp & Janka 2000; Mezzacappa et al. 2001; Liebendörfer et al. 2005). The exception being the $8.8 M_{\odot}$ progenitor of Nomoto & Hashimoto (1988) which resulted in an $\sim 3 \times 10^{49}$ erg neutrino-driven-wind explosion due to the rarefied stellar envelope (Kitaura, Janka & Hillebrandt 2006). Extension to axisymmetric calculations (2D) yielded similar outcomes over their 1D counterparts despite the inclusion of instabilities such as neutrino-driven convection and the standing accretion shock instability (SASI; Burrows et al. 1995; Janka & Mueller 1996; Buras et al. 2006a,b; Ott et al. 2008; Marek & Janka 2009).

It should be noted that while many of the current numerical experiments incorporate multidimensional hydrodynamics, performing 3D radiation is currently prohibitive computationally (Zhang et al. 2013). Many groups utilize multi group flux-limited diffusion (MGFLD) in the 1D ‘ray-by-ray’ transport approximation. This is a reasonable approach to core-collapse simulations both because of the limitation of current computational resources and because the results for multi-angle transport are similar to those for MGFLD except in the cases of extremely rapid rotation (Ott et al. 2008). Thus, it is unlikely that future simulations that incorporate 3D transport will yield fundamental differences over current state-of-the-art calculations in terms of the viability of neutrino mechanism.

Recently, a number of groups have published 3D core-collapse simulations with differing computational approaches and various levels of sophistication (Nordhaus et al. 2010b; Hanke et al. 2012,

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2013; Couch 2013a; Dolence et al. 2013; Janka 2013; Couch & O’Connor 2014; Dolence, Burrows & Zhang 2014; Mezzacappa et al. 2014; Takiwaki, Kotake & Suwa 2014). Some groups find that the extra degree of freedom available in 3D simulations makes it easier to achieve shock revival over their axisymmetric counterparts (Nordhaus et al. 2010b; Dolence et al. 2013). On the other hand, several groups have found the opposite; namely that explosions are harder to achieve in 3D than 2D (Hanke et al. 2012, 2013; Couch 2013a; Janka 2013; Couch & O’Connor 2014; Takiwaki et al. 2014). If that is the case, then it may well be that the delayed-neutrino mechanism categorically fails and alternative mechanisms should be investigated.

In one recent case, axisymmetric calculations of 12, 15, 20 and 25 M_{\odot} progenitors successfully revived the shock with explosion energy estimates of ~ 0.3 – 0.9 foe (Bruenn et al. 2013, 2014; Mezzacappa et al. 2014). Their energy is supplied primarily by an enthalpy flux. This is actually a wind, mainly along the imposed symmetry axis, i.e. a collimated wind. This wind is driven by the inflowing (accreted) gas. Winds were suggested to power CCSN in the past (e.g. Burrows & Goshy 1993; Burrows et al. 1995), but were found to have limited contribution to the explosion for a more massive star than $8.8 M_{\odot}$.

Many CCSNe, e.g. some recent Type Ic SNe (Roy et al. 2013; Takaki et al. 2013), explode with kinetic energy of $\gtrsim 10$ foe. Neutrino-based mechanisms cannot account for such energies even under favourable conditions. For example, Ugliano et al. (2012) performed a set of simulations where the energy was artificially scaled to that of SN 1987A, and found that even if neutrino explosions do work for some CCSNe, no explosions with kinetic energy of > 2 foe are achieved. This scaling was achieved by artificially setting the inner boundary luminosity to obtain an explosion with an energy equal to that of SN 1987A. The delayed-neutrino mechanism must be shown to produce robust explosions with canonical SN energies for a range of progenitors if it is to continue to be a contender in core-collapse theory. Despite decades of effort with the most sophisticated physics to date, no current simulation has produced a successful 10^{51} erg SN. It is this fact that leads us to argue that the delayed-neutrino mechanism has a generic character that prevents it from exploding the star with an observed energy of 1 foe.

The delicate and problematic nature of neutrino-driven mechanisms were already revealed with 1D simulations, such that even the most sophisticated neutrino transport calculations were unable to explode stars for progenitor masses $\gtrsim 12 M_{\odot}$ (e.g. Liebendörfer et al. 2001). Multidimensional effects were then seen as necessary for triggering an explosion. The most common multidimensional processes that have been studied as a rescue for the delayed-neutrino mechanism were neutrino-driven convection (e.g. Burrows et al. 1995) and hydrodynamic instabilities, such as the SASI (Blondin, Mezzacappa & DeMarino 2003). These axisymmetric (2D) simulations have shown mixed and contradicting results. Most do not get an ‘explosion’ at all, while others obtain explosions with very little energy, i.e. $\ll 1$ foe (e.g. Suwa et al. 2010; Suwa 2014). In most of these cases where an ‘explosion’ is claimed, it is actually only shock revival and not a typical explosion, as the energy is much too low to explain most observed CCSNe.

In the past few years, the regime of 3D flow structures has been explored in more detail (e.g. Nordhaus et al. 2010b). The simulations have not reached any consensus on the outcome. While some show that it is easier to revive the shock in 3D than in 2D (e.g. Nordhaus et al. 2010b; Dolence et al. 2013), others showed the opposite (e.g. Couch 2013b; Hanke et al. 2013; Couch & O’Connor 2014; Takiwaki et al. 2014). Even in 3D simulations that success-

fully revive the shock, the energy is significantly lower than 1 foe. Recently, turbulence from convective burning in the Si/O shell was shown to aid shock revival (Couch & Ott 2015; Mueller & Janka 2014).

A recent demonstration of outcome sensitivity on initial setting is the two 3D studies by Nakamura et al. (2014) and Mösta et al. (2014). Nakamura et al. (2014) find an explosion energy of ~ 1 foe for a case with a rapid core rotation. For a rotation velocity of 0.2 times that rapid rotation, the explosion energy was only ~ 0.1 foe. They did not include magnetic fields. Mösta et al. (2014) included very strong magnetic fields in the pre-collapse core as well as a very rapid rotation, about twice as large as the rapid rotation case of Nakamura et al. (2014). Mösta et al. (2014) obtained jets but did not manage to revive the stalled shock and did not obtain any explosion.

The structure of this paper is as follows. In Section 2, we expand upon the argument presented in Papish & Soker (2012a) that the delayed-neutrino mechanism cannot achieve canonical SN energies. We consider the limitation of the delayed-neutrino mechanism from another perspective in Section 3. In Section 4 we discuss the role of progenitor perturbations and why contradicting results are common among the groups simulating neutrino-based mechanisms, and in Section 5 we discuss the energy available from recombination of free nucleons. A discussion of the collimated wind obtained by Bruenn et al. (2014) and our summary are given in Section 6.

2 TIME-SCALE CONSIDERATIONS

We start with simple time-scale considerations during the revival of the shock in a spherically symmetric outflow. The ‘gain region’ of the delayed-neutrino mechanism, i.e. where neutrino heating outweighs neutrino cooling, typically occurs in the region $r \simeq 100$ – 400 km (Janka 2001).

For an explosion to be initiated, the advection time-scale τ_{adv} should be larger than the heating time-scale τ_{heat} . This advection time-scale is the time needed for material to cross the gain region during accretion. Most core-collapse simulations fail when this condition is not fulfilled. When this condition is met, the internal energy can increase until there is enough energy to unbind the material and an explosion is initiated. At this point, the total energy of the gas in the gain region is very close to zero. From this time, the net heating adds up to the positive explosion energy. After the gas reaches large radii, $\gtrsim 1000$ km, heating becomes inefficient. It is true that some gas expands at a lower velocity and it is closer to the centre. However, density decreases and so does the neutrino optical depth that decreases below its initial value, such that neutrino heating becomes even less efficient. Material near the neutrinosphere has, by definition, a large optical depth. It can in principle absorb energy and expand. But this process is a neutrino-driven wind, which is not part of the delayed-neutrino mechanism, and was found to have limited contribution to the explosion (e.g. Burrows & Goshy 1993; Burrows et al. 1995). The time from the start of acceleration to the end of efficient heating is marked t_{esc} . From simulations $t_{\text{esc}} \simeq 50$ ms (Marek & Janka 2009; Bruenn et al. 2013, 2014). In Section 3 we find a similar time from a simple analytical estimate.

In fig. 2 of Bruenn et al. (2013), the shock is starting to expand and an explosion is initiated at time $t \simeq 200$ ms. At this time, the total positive energy is close to zero (fig. 4 in Bruenn et al. 2013). At that time the shock is at a distance of $r_s \simeq 400$ km. This shows that during the time the shock moves from 200 to 400 km, the total energy increases from a negative value to about zero. We take the time of zero energy to be the starting point of positive energy

accumulation, and use it to estimate the explosion energy. In the simulations of Bruenn et al. (2013) at time $t = 300$ ms, the shock is already at a distance of $r_s \simeq 1000\text{--}1500$ km. Some material is closer to the centre, but its density is lower than that at earlier times, opacity is lower and heating is inefficient. We note again the long duration of energy increase in the work of Bruenn et al. (2013, 2014) and Mezzacappa et al. (2014), where energy increases linearly with time for over a second, a time when the shock is already at a distance of $r_s \simeq 10\,000$ km. This linear growth of the energy can be explained by a strong neutrino-driven wind from the proton-neutron star. In the new 3D case presented by Mezzacappa et al. (2014), the shock radius position is similar to their results of 1D simulations where no explosion has been obtained.

A similar dynamic can be seen in fig. 4 of the 2D simulation of Marek & Janka (2009), where at time $t = 524$ ms the shock is at a radius of $r_s \simeq 200$ km. The shock moves outwards to 400 km at $t = 610$ ms, but then at time $t = 650$ ms the shock radius decreases back to 200 km. This shows that at that time the energy is about zero and is not positive. The acceleration time can be inferred from fig. 6 where the average shock moves from 400 to 700 km during ~ 50 ms. In each direction, the acceleration time lasts for ~ 50 ms. However, as the acceleration occurs at different times at different directions, the behaviour of the average shock radius gives the impression that the acceleration phase is longer than 50 ms.

For a neutrinosphere at $r_v \simeq 50$ km (e.g. Couch & O'Connor 2014), the neutrino ‘optical depth’ from r to infinity is given by

$$\tau_v \simeq 0.1(r/100\text{ km})^{-3} \quad (1)$$

(Janka 2001), where the typical electron neutrino luminosity is $L_v \simeq L_{\bar{\nu}} \simeq 5 \times 10^{52} \text{ erg s}^{-1}$ (e.g. Mueller et al. 2012). Over all, if the interaction occurs near a radius r in the gain region, the energy that can be acquired by the expanding gas is

$$E_{\text{shell}} \simeq t_{\text{esc}} \tau L_v \simeq 0.25 \left(\frac{t_{\text{esc}}}{50\text{ ms}} \right) \left(\frac{L_v}{5 \times 10^{52} \text{ erg s}^{-1}} \right) \times \left(\frac{r}{100\text{ km}} \right)^{-3} \text{ foe.} \quad (2)$$

Using a more typical radius of ~ 200 km for the acceleration region reduces the total energy to 0.03 foe. Non-spherical flows that allow some simultaneous inflow–outflow structure might under favourable conditions be expected to increase the energy by a factor of few to $\sim 0.1\text{--}0.3$ foe. This is consistent with numerical simulation results of the delayed-neutrino mechanism summarized in Section 1. It is interesting to note that Bethe & Wilson (1985) found an explosion energy limit of 0.4 foe. This was based on their simulations and not on any physical reason why the neutrino mechanism fails.

3 ENERGY CONSIDERATIONS

We examine the situation by considering in more detail the acceleration from the delayed-neutrino mechanism. Consider a mass M_a that is accelerated and ejected by absorbing a fraction f of the neutrino energy. The mass starts at radius r_0 with zero energy. Namely the sum of internal and gravitational energy is zero. This is an optimistic assumption, as the internal energy itself also needs to be supplied by neutrinos. Neutrino losses can be absorbed into the parameter f . After an acceleration time t , the energy of the mass is $fL_v t$ and its velocity is

$$v = \frac{dr}{dt} \simeq \left(\frac{2fL_v t}{M_a} \right)^{1/2}. \quad (3)$$

Here we assume that most of the energy is transferred to kinetic energy. Initially, more energy can be stored as thermal energy. However, not much thermal energy can be stored after the gas energy becomes positive, as it starts to accelerate outwards and thermal energy is converted to kinetic energy on a dynamical time-scale. The thermal energy acts to overcome gravity. We calculate here the extra energy that goes to gas outward motion.

Let the acceleration be effective to radius r_a at time t_a . Integrating over time gives

$$r_a - r_0 \simeq \frac{2}{3} \left(\frac{2fL_v}{M_a} \right)^{1/2} t_a^{3/2}, \quad (4)$$

or

$$\begin{aligned} t_a &\simeq \left(\frac{9}{8} \right)^{1/3} (r_a - r_0)^{2/3} \left(\frac{M_a}{fL_v} \right)^{1/3} \\ &= 0.05 \left(\frac{r_a - r_0}{500\text{ km}} \right)^{2/3} \left(\frac{M_a}{0.1 M_\odot} \right)^{1/3} \left(\frac{L_v}{5 \times 10^{52} \text{ erg s}^{-1}} \right)^{-1/3} \\ &\quad \times \left(\frac{f}{0.1} \right)^{-1/3} \text{ s.} \end{aligned} \quad (5)$$

A similar acceleration time is estimated from numerical results as we discussed in Section 2, where this time is marked t_{esc} . Under these assumptions, the energy of the ejected mass is

$$\begin{aligned} E_a &\simeq t_a f L_v \simeq 0.24 \left(\frac{r_a - r_0}{500\text{ km}} \right)^{2/3} \left(\frac{M_a}{0.1 M_\odot} \right)^{1/3} \\ &\quad \times \left(\frac{L_v}{5 \times 10^{52} \text{ erg s}^{-1}} \right)^{2/3} \left(\frac{f}{0.1} \right)^{2/3} \text{ foe.} \end{aligned} \quad (6)$$

In these calculations, we assumed a constant neutrino luminosity. As the neutrino luminosity decreases with time (e.g. Fischer et al. 2012), the term fL_v , in equation (6), actually overestimates the available energy. More typical values for acceleration over ~ 500 km are $f < 0.1$ due to the low neutrino opacity (equation 1) and lower accelerated mass. These values give $E_a < 0.2$ foe as in equation (2).

4 THE ALMOST UNBOUND STALLED SHOCK

The energy of the immediately post-shocked gas falling from thousands of km to hundreds of km is close to zero before there is much neutrino cooling. Whether the shocked gas falls or expands is a question of whether a small amount of energy is added to revive the shock. When there are departures from spherical symmetry, like the perturbations introduced by Couch & Ott (2013) or instabilities in the post-shock region, in some areas the extra energy comes at the expense of other areas. For example, a vortex can add a positive velocity in the region of the flow where the flow goes out. Even if the shock is revived, the energy limitations given in Sections 2 and 3 apply. The SASI itself is a manifestation of the process where one region of the stalled shock can go out in expense of other regions. The extra energy from neutrino heating can even revive the entire sphere. However, the energy gained by neutrino heating is limited.

A recent attempt to revive the stalled shock is that of Couch & Ott (2013), who introduced perturbations to the Si/O layers and found them to enable shock revival under certain conditions. What Couch & Ott (2013) term a successful explosion is actually a revival of the stalled shock. They did not obtain the desired ~ 1 foe explosion. As with many other simulations, small changes in the initial conditions determine whether shock revival occurs or not. For example, Couch & Ott (2013) find shock revival when their neutrino heat factor is

1.02, but not when it is 1. They present their average shock position until it reaches a radius of 430 km at $t = 0.32$ s. Examining their successful revival run presented in their fig. 3, we find the average shock outward velocity in the last part they show, 370–430 km, to be $(v_{\text{shock}}) = 8000 \text{ km s}^{-1}$. This is less than 0.3 times the escape velocity at that radius. The shock does not seem to accelerate in the last 50 km. Within $\Delta t \simeq 0.04$ s, the shock will reach a radius of about 700 km, where no more energy gain is possible (Janka 2001). At 400 km the neutrino optical depth is very small, $\tau < 0.1$. Indeed, at an average shock radius of 350 km, the heating efficiency in their simulation η , defined as net heating rate divided by $L_{\nu_e} + L_{\bar{\nu}_e}$, drops below 0.1. This implies that the gained energy will be very small, $\Delta E < \tau L_{\nu} \Delta t < 0.2$ foe. We therefore estimate that even the perturbations introduced by Couch & Ott (2013) will not bring the explosion, if occurs, close to 1 foe.

Let us quantify the statement of energy close to zero. We can make the following estimations based on the models of Woosley, Heger & Weaver (2002) of massive stars prior to the collapse. The gas at 2000 km has a specific gravitational energy of $e_{G0} = -10^{18} \text{ erg g}^{-1}$ and a specific internal energy of $e_{i0} = 5.5 \times 10^{17} \text{ erg g}^{-1}$. After mass-loss to neutrinos from the core, the inner mass reduces by ~ 10 per cent. However, by that time the shell that starts at $\text{few} \times 1000$ km has been accelerated inwards. So we take the total specific energy to be the pre-collapse energy. As an example, we take the stalled shock to be at $r_s = 200$ km. When reaching $r_s = 200$ km, the specific total energy $e_t = e_{i0} + e_{G0}$ stays the same. The specific gravitational energy is $e_{Gs} \simeq 10e_{G0} = -10^{19} \text{ erg g}^{-1}$, and the specific internal (thermal + kinetic + nuclear) energy is $e_{Is} = e_t - 10e_{G0} \simeq 9.5 \times 10^{18} \text{ erg g}^{-1}$. The net specific energy relative to gravitational energy in this demonstrative example is

$$\xi_s \equiv \left| \frac{e_{Is}}{e_{Gs}} \right| \simeq 0.95. \quad (7)$$

The mass is very close to be unbound. Small amount of net heating can revive the shock. For a typical mass in the gain region of $M_{\text{gain}} \lesssim 0.05 M_{\odot}$ (e.g. Couch 2013b), an extra energy of $\Delta E = 5 \times 10^{49} \text{ erg} = 0.05$ foe will revive the shock.

5 ENERGY AVAILABLE FROM RECOMBINATION

Adding nuclear energy of free nucleons does not change the above property of an almost unbound stalled shock and the conclusion of low ‘explosion’ energy. Consider the scenario where disintegration of nuclei forms free nucleons beyond the stalled shock, and the available nuclear energy is reused later after the free nucleons are accelerated outwards by neutrinos (Janka et al. 2012). When the nucleons recombine to form heavy nuclei, an energy of up to 9 MeV per nucleon can in principle be used to explode the star (Janka et al. 2012). A mass of $0.06 M_{\odot}$ in the gain region can then release in principle $\sim 10^{51} \text{ erg}$ (Scheck et al. 2006).

However, the recombination of free nucleons to alpha particles, a process that uses 7 MeV from the 9 MeV available in forming silicon, starts when the reviving post-shock gas reaches $r \sim 250$ km (Fernández & Thompson 2009). The energy released by recombination accelerates the material (Fernández & Thompson 2009), which results in a shorter acceleration time than given in equation (5). This further lowers the energy that can be supplied by neutrinos below that given in equations (2) and (6).

The energy available from recombination is limited as well. From fig. 5 of Fernández & Thompson (2009), we find the total fraction

of α particles in the gas inside the shock when the shock radius is 500 km to be $X_{\alpha} \lesssim 0.5$; the fraction just behind the shock front at 500 km is $X_{\alpha} \simeq 0.9$. In the results of Fernández & Thompson (2009), the fraction of α particles increases as the shock moves outwards. For this fraction, the average energy available from recombination is 5 MeV per nucleon (Janka et al. 2012). However, the shock is only at 500 km and a large fraction of the mass is much deeper. As the shock expands further, the fraction of α particles will increase and the available energy will decrease. Taking a mass in the gain region of $M_{\text{gain}} \lesssim 0.05 M_{\odot}$ (e.g. Couch 2013b), we find the ‘explosion’ energy to be $E_{\text{nuc}} \lesssim 0.5$ foe. In some 2D simulations, the mass in the gain layer is $\gtrsim 0.05 M_{\odot}$, but in 3D simulations the gain layer has lower mass than in 2D simulations (e.g. Couch 2013a). Over all, the available energy without neutrino winds or jets is < 0.5 foe. This value is an upper limit and consistent with many of the simulations summarized in Section 1 that achieve much lower energies or do not revive the shock at all.

It should be emphasized that the recombination is not a new energy source, as the thermal energy of the shocked gas is used to disintegrate the nuclei. The recombination is the re-usage of this energy. The extra energy must come from neutrinos that lift the free nucleons to larger radii. The total available energy from recombination is proportional to the mass of the free nucleons that are lifted from small $r \lesssim 150$ km to large radii $r \gtrsim 500$ km. However, the amount of mass that can be accumulated at small radii is limited because if the density is too high, then cooling overcomes neutrino heating, and the shock will not be revived.

Yamamoto et al. (2013) performed 1D and 2D simulations of shock revival and examined explosion energy including recombination and shock nuclear burning. They tuned the neutrino luminosity to a critical value that gives successful explosions. Their successful runs have shock relaunch times of 0.3–0.4 s in 2D flows. The explosion energy in these runs is in the range of 0.6–1.5 foe. We note the following regarding their tuned calculations.

(1) Yamamoto et al. (2013) assume that neutrino heating alone revives the stalled shock. Then they can use the entire recombination energy to explode the rest of the star. The more realistic calculations of Fernández & Thompson (2009) show that at least half the recombination energy is required to help revive the shock.

(2) The above assumption implies the need for high neutrino luminosity. Indeed, in Yamamoto et al. (2013) successful 2D runs the required critical neutrino luminosities are $L_{\nu_e, c} = L_{\bar{\nu}_e, c} = 4.8 \times 10^{52}$ and $4.5 \times 10^{52} \text{ erg s}^{-1}$ for shock relaunching times of 0.3 and 0.4 s, respectively. These neutrino luminosities are ~ 50 per cent higher than what most realistic numerical simulations find, e.g. Fischer et al. (2012), and ~ 30 per cent higher than the neutrino luminosities obtained by Mueller et al. (2012) who included general relativistic effects. Interestingly, Mueller et al. (2012) find for their $11 M_{\odot}$ model that recombination of nucleons and α particles in the ejecta would provide an additional energy of $E_{\text{rec}} \simeq 0.02$ foe. For their $15 M_{\odot}$ model, they argue that burning in the shock will add of the order of 0.1–0.2 foe or more.

(3) The contribution of nuclear and recombination energies to the diagnostic explosion energy of Yamamoto et al. (2013) is very similar to the contribution of neutrino heating.

Based on these points we can use a more realistic value of neutrino heating, $E_{\nu} < 0.2$ foe, and conclude that the combined explosion energy in realistic simulations will be $E_{\text{exp}} < 0.5$ foe. Again we reach the conclusion that including recombination energy will at most bring the explosion energy to $E_{\text{exp}} < 0.5$ foe. Although close to the canonical 1 foe value, one must keep in mind that this

value is obtained with very favourable conditions, and in scaled, rather than realistic, simulations. In more realistic simulations, the recombination energy is found to be $E_{\text{rec}} \lesssim 0.2$ foe, e.g. Mueller et al. (2012).

6 DISCUSSION AND SUMMARY

Using simple estimates of a spherically symmetric mass ejection by neutrino flux in CCSNe, we found that in the delayed-neutrino mechanism (Bethe & Wilson 1985), where the main energy source of the explosion is due to neutrino heating in the gain region, the explosion energy is limited to $E_{\text{exp}} \lesssim 0.5$ foe, with a more likely limit of 0.3 foe (equations 2 and 6). This falls short of what is required in most CCSNe.

Although our simple analytical estimates are limited to spherically symmetric outflows, they none the less catch the essence of the delayed-neutrino mechanism. In a non-spherical flow, instabilities, such as neutrino-driven convection and the SASI, play a major role (e.g. Hanke et al. 2013). Such instabilities allow inflow and outflow to occur simultaneously. Still, recent and highly sophisticated 3D simulations with enough details to resolve such instabilities do not obtain enough energy to revive the stalled shock (e.g. Janka 2013). The energy that can be used from the neutrino flux might, under favourite conditions, revive the stalled shock, but cannot lead to explosions with energies of $E_{\text{exp}} \gtrsim 0.3$ foe.

Our conclusion holds as long as no substantially new ingredient is added to the delayed-neutrino mechanism. Such an ingredient can be a strong wind, as was applied by artificial energy deposition by Nordhaus et al. (2010a, 2012). In their 2.5D simulations, Scheck et al. (2006) achieved explosion that was mainly driven by a continuous wind. The problem we see with winds is that they are less efficient than jets. Indeed, in order to obtain an explosion, the winds in the simulations of Scheck et al. (2006) had to be massive. For that, in cases where they obtained energetic enough explosions, the final mass of the NS was low ($M_{\text{NS}} < 1.3 M_{\odot}$). Such a wind must be active while accretion takes place; the accretion is required to supply the energy (Marek & Janka 2009).

With the severe problems encountered by neutrino heating, research groups have turned to study dynamical processes. Couch & Ott (2013, 2015) and Mueller & Janka (2014) argued that the effective turbulent ram pressure exerted on the stalled shock allows shock revival with less neutrino heating than 1D models. However, Abdikamalov et al. (2014) found that increasing the numerical resolution allows cascade of turbulent energy to smaller scales, and the shock revival becomes harder to achieve at high numerical resolutions.

Another dynamical process is a collimated wind blown by the newly formed NS. Bruenn et al. (2014) performed 2D simulations up to 1.4 s post-bounce, and obtained an explosion energy of 0.3–0.9 foe, depending on the stellar model (initial mass without rotation). They find the main energy source to be what they term an ‘enthalpy flux’. This is actually a wind, mainly along the imposed symmetry axis, i.e. a collimated wind. This wind is driven by the inflowing (accreted) gas. At some instant, their results show jet-like outflows along the symmetry axis. It seems that the collimated wind is induced by the numerical grid. Contrast that to their corresponding 3D simulations (Mezzacappa et al. 2014) which show no such explosion. Mezzacappa et al. (2014) present one new result of a 3D run for their $15 M_{\odot}$ model at $t = 267$ ms post-bounce. We estimate the average shock radius at that time to be ~ 220 km. This is very similar to their 1D results (Bruenn et al. 2013), where the shock radius is much smaller than that in their 2D simulations, and where no

explosions occur. None the less, the results of Bruenn et al. (2014) show the great potential of an inflow–outflow mechanism in exploding CCSNe. An inflow–outflow situation with collimated outflows over a relatively long time naturally occurs with jets launched by accretion discs, without the numerically induced symmetry axis in 2D grids.

For the above, the lack of persisting success, and possibly failure, of the delayed-neutrino mechanism calls for a paradigm shift. As well, the rich variety of CCSN properties (e.g. Arcavi et al. 2012) further emphasizes the need to study alternative models for CCSN explosions, some of which are based on jet-driven explosions (Janka 2012). In CCSN simulations, jets have been shown to be launched when the pre-collapsing core possesses both a rapid rotation and a very strong magnetic field (e.g. LeBlanc & Wilson 1970; Bisnovaty-Kogan, Popov & Samokhin 1976; Meier et al. 1976; Khokhlov et al. 1999; Höflich, Khokhlov & Wang 2001; MacFadyen, Woosley & Heger 2001; Woosley & Janka 2005; Burrows et al. 2007; Couch, Wheeler & Milosavljević 2009; Couch et al. 2011; Takiwaki & Kotake 2011; Lazzati et al. 2012). However, these jets do not explode the core via a feedback mechanism, such that they too often give extreme cases as gamma-ray bursts, or they fail to explode the star, e.g. Mösta et al. (2014). Recent observations (e.g. Lopez et al. 2013; Milisavljevic et al. 2013) suggest that jets might play a role in at least some CCSNe. Another motivation to consider jet-driven explosion mechanisms is that jets might supply the site for the r-process (Papish & Soker 2012b; Winteler et al. 2012). The question is whether the accreted mass possesses sufficient specific angular momentum to form an accretion disc. Persistent accretion disc requires the pre-collapsing core to rotate fast, as in the magnetohydrodynamics class of models (e.g. LeBlanc & Wilson 1970; Bisnovaty-Kogan et al. 1976; Meier et al. 1976; Khokhlov et al. 1999; Höflich et al. 2001; MacFadyen et al. 2001; Woosley & Janka 2005; Burrows et al. 2007; Couch et al. 2009, 2011; Lazzati et al. 2012). Most massive stars reach the core-collapse phase with a too slow core rotation for the magnetorotational mechanism to be significant.

One alternative to the delayed-neutrino mechanism which overcomes the angular momentum barrier is the so-called jittering-jet mechanism of Papish & Soker (2011). The jittering-jet mechanism overcomes the requirement for rapid core rotation, and was introduced as a mechanism to explode all CCSNe (Papish & Soker 2011, 2012b, 2014). The angular momentum source is the convective regions in the core (Gilkis & Soker 2014), and/or instabilities in the shocked region of the collapsing core. Blondin & Mezzacappa (2007), Fernández (2010) and Rantsiou et al. (2011) suggested that the source of the angular momentum of pulsars is the spiral mode of the SASI. In the jittering-jet mechanism, there is no need to revive the accretion shock, and it is a mechanism based on a negative feedback cycle. As long as the core was not exploded, the accretion continues. After an energy several times the core binding energy is deposited to the core by the jets, the star explodes. This energy amounts to ~ 1 foe. If the feedback is less efficient, more accretion is required to accumulate the required energy. If the efficiency is very low, the accreted mass on to the NS brings it to collapse to a black hole and launch relativistic jets. Namely, in general, the less efficient the feedback mechanism is, the more violent the explosion is (Gilkis & Soker 2014).

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