NEUTRINO-DRIVEN EXPLOSION OF A 20 SOLAR-MASS STAR IN THREE DIMENSIONS ENABLED BY STRANGE-QUARK CONTRIBUTIONS TO NEUTRINO–NUCLEON SCATTERING

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ABSTRACT

Interactions with neutrons and protons play a crucial role for the neutrino opacity of matter in the supernova core. Their current implementation in many simulation codes, however, is rather schematic and ignores not only modifications for the correlated nuclear medium of the nascent neutron star, but also free-space corrections from nucleon recoil, weak magnetism, or strange quarks, which can easily add up to changes of several 10% for neutrino energies in the spectral peak. In the Garching supernova simulations with the PROMETHEUS-VERTEX code, such sophistications have been included for a long time except for the strange-quark contributions to the nucleon spin, which affect neutral-current neutrino scattering. We demonstrate on the basis of a 20 \( M_\odot \) progenitor star that a moderate strangeness-dependent contribution of \( g^s_\nu = -0.2 \) to the axial-vector coupling constant \( g_\nu \approx 1.26 \) can turn an unsuccessful three-dimensional (3D) model into a successful explosion. Such a modification is in the direction of current experimental results and reduces the neutral-current scattering opacity of neutrons, which dominate in the medium around and above the neutrinosphere. This leads to increased luminosities and mean energies of all neutrino species and strengthens the neutrino-energy deposition in the heating layer. Higher nonradial kinetic energy in the gain layer signals enhanced buoyancy activity that enables the onset of the explosion at \( \sim 300 \) ms after bounce, in contrast to the model with vanishing strangeness contributions to neutrino–nucleon scattering. Our results demonstrate the close proximity to explosion of the previously published, unsuccessful 3D models of the Garching group.

Key words: hydrodynamics -- instabilities -- neutrinos -- supernovae: general

Supporting material: 3D figure, animation

1. INTRODUCTION

According to the standard paradigm of the explosion mechanism of core-collapse supernovae (SNe), neutrino heating above the gain radius initiates the outward acceleration of the stalled bounce shock and provides a major fraction of the energy that unbinds the explosion ejecta and powers the SN blast wave (e.g., Colgate & White 1966; Bethe & Wilson 1985; Bethe 1990; Janka 2012; Burrows 2013). Except for stars near the low-mass end of SN progenitors with oxygen-neon-magnesium cores or small iron cores and a steep density gradient of the overlying shells, successful explosions cannot be obtained in spherical symmetry (1D) by state-of-the-art simulations. For more massive progenitors the success of the neutrino-driven mechanism requires the support by multidimensional hydrodynamic flows in the postshock layer associated with convective overturn and the standing accretion shock instability (SASI; Blondin et al. 2003). Such flows increase the neutrino-energy deposition and create buoyancy and turbulent pressure, thus reducing the critical threshold for the neutrino luminosity to trigger the explosive runaway of the shock (e.g., Herant et al. 1994; Burrows et al. 1995; Janka & Müller 1996; Murphy & Burrows 2008; Nordhaus et al. 2010; Hanke et al. 2012; Dolenčec et al. 2013; Murphy et al. 2013). Breaking spherical symmetry also allows for simultaneous mass accretion and outflow from the proto-neutron star (PNS) after shock revival, which steepens the rise of the explosion energy (Marek & Janka 2009).

The first self-consistent three-dimensional (3D) stellar core-collapse simulations with successful explosions were performed by Fryer & Warren (2002, 2004), using smoothed-particle hydrodynamics and a simple, gray diffusion description of neutrino transport, which favored rapid explosions in 2D and 3D. Recently, the Garching group has obtained a 3D explosion for a 9.6 \( M_\odot \) star with sophisticated, state-of-the-art ray-by-ray-plus (RbR+) multi-group transport and found that 3D hydrodynamic flows enhance the explosion energy compared to the 2D (axisymmetric) case (Melson et al. 2015). However, previous 3D simulations of this group for 11.2, 20, and 27 \( M_\odot \) progenitors (covering \( \sim 400–550 \) ms after bounce) did not produce explosions although corresponding 2D models exploded (Hanke et al. 2013; Tamborra et al. 2013, 2014a, 2014b). This is consistent with studies based on parametrized neutrino heating (e.g., Hanke et al. 2012; Couch 2013; Couch & O’Connor 2014; Couch & Ott 2015) and self-consistent simulations with approximate neutrino transport (Takiwaki et al. 2012, 2014), which revealed that the onset of explosion is less favored or delayed in 3D compared to 2D, contradicting claims by Nordhaus et al. (2010) and, though moderated, by Burrows et al. (2012) and Dolenčec et al. (2013). Also Lentz et al. (2015), using high-fidelity RbR multi-group neutrino diffusion, highlight a successful 15 \( M_\odot \) explosion that sets in significantly later than the corresponding 2D model. Large-scale deformation modes seem to have the strongest supportive effect on the explosion but are weakened by the forward turbulent energy cascade to small-scale structures in 3D, which is opposite to the 2D case.
While the resolution feasible in current full-scale 3D SN models is still insufficient to satisfactorily represent postshock turbulence (Abdikamalov et al. 2014; Radice et al. 2015), the results reported above suggest that important physics might still be missing in the models. One of the aspects to be scrutinized are the pre-collapse initial conditions, which result from 1D stellar evolution modeling. Couch & Ott (2013), Couch et al. (2015), and Müller & Janka (2015) indeed confirmed speculations that large-amplitude perturbations of low-order modes in the convective shell-burning layers (e.g., Arnett & Meakin 2011, and references therein) might facilitate the development of explosions. Further progress will require 3D modeling of the final stages of stellar evolution.

Here we demonstrate that remaining uncertainties in the neutrino opacities, in particular the neutrino-strange-quark contribution of neutral-current scatterings. We show that a moderate isoscalar strange-quark contribution to the nucleon spin in their effect on weak collapse simulations. As an example we consider possible strange-

3. STRANGENESS CONTRIBUTIONS TO NEUTRINO–NUCLEON SCATTERING

The lowest-order differential neutrino–nucleon scattering cross section reads

\[ \frac{d\sigma_0}{d\Omega} = \frac{G_F^2 \epsilon^2}{4\pi^2} \left[ c_a^2 (1 + \cos \theta) + c_a^2 (3 - \cos \theta) \right]. \tag{1} \]

with \( \epsilon \) being the incoming neutrino energy, \( \theta \) the scattering angle, \( G_F \) Fermi’s constant, and \( c_a \) and \( c_a \) vector and axial-vector coupling constants, respectively. The latter are

\[ c_a = \frac{1}{2} - 2 \sin^2 \theta_W \approx 0.35, \quad c_a = g_a/2 \approx 0.63 \quad \text{for } \nu p \rightarrow \nu p \]

and

\[ c_a = -\frac{1}{2}, \quad c_a = -g_a/2 \approx -0.63 \quad \text{for } \nu m \rightarrow \nu m \quad \text{with } g_a \approx 1.26 \quad \text{and } \sin^2 \theta_W \approx 0.2325. \]

For iso-energetic scattering \((\epsilon' = \epsilon)\), Equation (1) yields the total transport cross section

\[ \sigma_0 = \int_{4\pi} d\Omega \frac{d\sigma_0}{d\Omega} (1 - \cos \theta) = \frac{2G_F^2 \epsilon^2}{3\pi} (c_a^2 + 5c_a^2). \tag{2} \]

While in our SN simulations corrections due to nuclear thermal motions and recoil, weak magnetism, and nucleon correlations at high densities are taken into account (Rampp & Janka 2002; Buras et al. 2006), Equations (1), (2) provide good estimates. Strange-quark contributions to the nucleon spin modify \( c_a \) according to

\[ c_a = \frac{1}{2} (\pm g_a - g_a^s), \tag{3} \]

where the plus sign is for \( \nu p \) and the minus sign for \( \nu m \) scattering (see, e.g., Horowitz 2002; Langanke & Martínez-Pinedo 2003). Since \( g_a^s < 0 \), the cross section for \( \nu p \)-scattering is increased and for \( \nu m \)-scattering decreased.

Employing Equation (2) with \( g_a^s = -0.2, \) Horowitz (2002) estimates 15%, 21%, 23% reduction of the neutral-current opacity for a neutron–proton mixture with electron fractions \( Y_e = 0.2, 0.1, 0.05 \), which are typical values for the layer between neutrinosphere (at density \( \rho \sim 10^{11} \text{ g cm}^{-3} \)) and \( \rho \sim 10^{13} \text{ g cm}^{-3} \) for hundreds of milliseconds after bounce. Since strangeness does not affect charged-current interactions and NS matter is neutron-rich, the reduced scattering opacity allows mainly heavy-lepton neutrinos \((\nu_e \equiv \nu_x, \bar{\nu}_x, \nu_s, \bar{\nu}_s)\) to leave the hot accretion mantle of the PNS more easily. This was found to enhance the expansion of the stalled SN shock in 1D models, although not enough for successful shock revival (Liebendörfer et al. 2002; Langanke & Martínez-Pinedo 2003). However, below we will show that the situation can be fundamentally different in 3D simulations.

4. RESULTS

We compare 2D and 3D core-collapse simulations of the 20 \( M_\odot \) star with strangeness corrections in neutrino–nucleon scatterings, using \( g_a^s = -0.2 \) (models 2Ds, 3Ds), to corresponding simulations without strange-quark effects \((g_a^s = 0); \) models 2Dn, 3Dn) as in all SN simulations of the Garching group so far. To explore “extreme” effects, our choice of \( g_a^s \) is by its absolute value somewhat bigger than theoretical and experimental determinations of \( g_a^s \sim 0.1 \) (Ellis & Kargler 1997; Airapetian et al. 2007; Alexakhin et al. 2007).

4.1. Dynamical Evolution toward Explosion

The sequence of 3D images in Figure 1 shows the postbounce evolution of the exploding model 3Ds; the entropy cuts in Figure 2 demonstrate the differences to the unsuccessful model 3Dn. In both cases the dynamics of the accretion layer are strongly SASI-dominated. In model 3Ds there is no indication of postshock convection but SASI sloshing motions first appear at post-bounce (p.b.) time \( t_{pb} \sim 120 \text{ ms} \).
These reach full strength around 180 ms p.b. and continue in varying directions until \(\sim280\) ms. Only later on convective overturn takes over as the dominant non-radial instability. In model 3Dn moderate buoyancy activity is visible from \(100-180\) ms before SASI becomes dominant, too. Early convection is enabled by a slightly larger radius of the stalled shock in 3Dn as a consequence of a slightly lower mass-accretion rate \(\dot{M}\) (Figure 3). This difference improves the growth conditions for postshock convection. In 3Dn an erroneous change of the transition between low-density and high-density EOS caused a delay of the core collapse. Therefore the mass-accretion rate until \(t \sim 150\) ms is slightly reduced and the Si/Si+O interface arrives at the shock \(\sim 15\) ms later. These early differences are inessential for our discussion because neutrino heating creates favorable conditions for the explosion of model 3Ds only later than about 300 ms.

From \(t_{pb} \sim 170\) ms on, 3Ds exhibits clearly larger SASI amplitudes and higher postshock entropies, which increases the maximum and average shock radii (Figures 2, 3). This model also shows larger non-radial kinetic energies in the gain layer, 

\[
E_{gain}^{\text{kin,\theta,\phi}} = \int_{R_{\text{gain}}}^{R_{\text{shock}}} dV \frac{1}{2} \rho \left(v_\theta^2 + v_\phi^2\right)
\]

(Figure 3), except during \(t_{pb} \sim 200-300\) ms, when a powerful spiral SASI mode develops in model 3Dn but not in 3Ds, albeit without pushing the shock sufficiently far out for revival (in conflict with recent results by Fernández 2015).

At \(t_{pb} \gtrsim 300\) ms, roughly 50 ms after the Si/Si+O interface has fallen through the shock and the shock has expanded to \(\sim 150\) km, the evolutions of models 3Dn and 3Ds separate. While in 3Dn the average shock radius, \(\langle R_{\text{shock}}\rangle\), retreats again...
and conditions become unfavorable for an explosion, model 3Ds exhibits positive trends in all explosion-diagnostic parameters (e.g., $R_{\text{shock}}$, $E_{\text{kin}}$, $Q_{\text{gain}}$, $M_{\text{gain}}$). Continuous shock expansion signals runaway and finally outward acceleration sets in at $t_{\text{shock}} \gtrsim 360$ ms, at which time the recombination of free nucleons to $\alpha$-particles in the largest plumes begins to release energy and the instantaneous “diagnostic energy,”

\[ E_{\text{exp}} = \int_{v_e > 0, \text{postshock}} dV \rho \epsilon_{\text{tot}}, \]  

starts to rise steeply, correlated with a fast growth of the mass in the gain layer (Figure 3). In Equation (5) the volume integration is performed over the postshock region where the total specific energy,

\[ \epsilon_{\text{tot}} = e + \frac{1}{2} |v|^2 + \Phi + \left[ e_{\text{bind}} \left( ^{56}\text{Fe} \right) - e_{\text{bind}} \right], \]

is positive, with $e$, $\frac{1}{2} |v|^2$, and $\Phi$ being the specific internal, kinetic, and (Newtonian) gravitational energies. The bracketed term expresses the difference between the specific nuclear binding energy when all nucleons are finally recombined to iron-group nuclei compared to the nuclear composition at a given time. It therefore accounts for the maximum release of nuclear energy and corresponds to an upper limit of $E_{\text{exp}}$.

Both corresponding 2D models, 2Dn and 2Ds, also explode after the Si/Si+O interface has passed the shock, but in 2Ds the outward shock acceleration and rise of $E_{\text{exp}}$ sets in $\sim 50$ ms earlier (Figure 3). Strangeness corrections therefore create more favorable explosion conditions also in the 2D case, although their influence is modest in successful models, similar to their small effect in the 1D case, which is far away from explosion (Liebendörfer et al. 2002).

### 4.2. Strangeness Corrections and Explosion

The impact of strangeness effects on the neutrino emission and the explosion is displayed in Figure 4. Model 3Ds consistently exhibits higher luminosities and mean energies of the emission for all neutrino species, and, consequently, a higher neutrino-energy deposition rate in the gain layer, $Q_{\text{gain}}$, a higher heating efficiency,

\[ \eta_{\text{heat}} = \frac{Q_{\text{gain}}}{L_{\nu_e} + L_{\nu_\mu}}, \]
a smaller gain radius (Figure 3), and a shorter heating timescale,

\[ \tau_{\text{heat}} = \frac{|E_{\text{gain}}|}{Q_{\text{gain}}} \]

(8)

with \( E_{\text{gain}} = \int_{R_{\text{gain}}} dV \rho (e + \frac{1}{2} |\mathbf{v}|^2 + \Phi) \) being the binding energy of the gain layer. Since the effective timescale of mass advection through the gain layer,

\[ \tau_{\text{adv}} = \frac{M_{\text{gain}}}{M} \]

(9)

(where \( M > 0 \)), which measures the average exposure time of matter to neutrino heating, is very similar in models 3Ds and 3Dn, the smaller \( \tau_{\text{heat}} \) in 3Ds also leads to a higher timescale ratio \( \tau_{\text{adv}}/\tau_{\text{heat}} \). The ratio \( \tau_{\text{adv}}/\tau_{\text{heat}} \) exceeds the critical value of unity shortly before the SN shock in 3Ds begins its runaway expansion.

The mean energies of the radiated neutrinos in model 3Ds are up to \( \sim 1 \text{ MeV} \) higher and the luminosities of \( \nu_e \) and \( \bar{\nu}_e \) by up to \( \sim 10\%-15\% \), whereas the \( \nu_x \)-luminosities rise by up to \( \sim 30\% \). The increase of the total neutrino luminosity is more than \( 6 \times 10^{52} \text{ erg s}^{-1} \) at maximum, which mainly comes from layers below the \( \nu_e \)-sphere, because the neutrino-loss rate \( Q_{\text{cool}} \) between the location of this sphere (at \( \sim 10^{11} \text{ g cm}^{-3} \)) and the gain radius differs between models 3Ds and 3Dn by at most \( \sim 10^{52} \text{ erg s}^{-1} \) (Figure 4). Note that at \( t_{\text{pb}} \gtrsim 300 \text{ ms} \) the relative differences of the neutrino properties of models 3Ds and 3Dn decrease and even change sign, because the former explodes whereas the latter continues to collapse and to accrete mass onto the PNS at a higher rate.
For $Y_e = 0.1-0.05$, strangeness effects in the neutrino–nucleon interactions reduce the effective opacity $\kappa_{\text{eff}} = \sqrt{\kappa_{\text{abs}}(\kappa_{\text{abs}} + \kappa_{\text{scatt}})}$ for $\nu_e$ only by 2%-3% and for $\bar{\nu}_e$ by 8%-10%. A considerable part of the observed luminosity enhancement of $\nu_e$ and $\bar{\nu}_e$ is therefore caused indirectly by a stronger contraction of the PNS in response to the larger energy loss through $x_n$ emission. With the smaller PNS radius (Figure 3) and steeper density profile, the neutrinospheres of $\nu_e$ and $\bar{\nu}_e$ move inward to higher temperatures.

The reduction of the weak neutral-current scattering by strange-quark contributions to the nucleon spin therefore enhances the neutrino luminosities and mean energies directly and indirectly and thus strengthens the neutrino heating in the gain layer. This amplifies buoyancy and turbulent mass motions behind the shock, which is signaled by higher non-radial kinetic energy in model 3Ds at $t_{pb} \gtrsim 300$ ms (Figure 3), thus fostering the explosion of this model in contrast to 3Dn.

### 4.3. Further Development of Explosion Energy

At the end of the simulation at 530 ms p.b., the PNS has a baryonic mass of $1.91 M_\odot$ and the shock has expanded to an average radius of $\langle R_{\text{shock}} \rangle \approx 1000$ km. The diagnostic energy, $E_{\text{exp}}$, has reached $0.2 \times 10^{51}$ erg and rises steeply and linearly with a rate of $\dot{E}_{\text{exp}} \approx 1.2 \times 10^{51}$ erg s$^{-1}$. This growth rate can be understood by the ejection of a continuous outflow of freshly neutrino-heated matter, $M_{\text{shock}} > 0 \approx 0.08-0.1 M_\odot$ s$^{-1}$ (Figure 5; the subscript indicates positive radial velocity), which is mainly fed by the shock-accreted matter that is channeled to the gain radius in persistent accretion downdrafts (Figure 2) with mass-inflow rates up to $\sim 0.3 M_\odot$ s$^{-1}$ (Figure 5). The mass outflow does not only absorb neutrino energy, it also releases nuclear binding energy from the recombinantion of neutrons and protons to $\alpha$-particles and heavy nuclei. Since neutrino heating roughly neutralizes the gravitational binding energy of the matter (Janka 2001; Marek & Janka 2009), the growth of $E_{\text{exp}}$...
can be estimated in terms of $M_{r>0}$ and an average nuclear recombination energy per nucleon, $\epsilon_{\text{nuc}}$, as

$$
E_{\text{exp}} \approx M_{r>0} \frac{\epsilon_{\text{nuc}}}{m_B} \approx 1.2 \times 10^{51} \left( \frac{M_{r>0}}{0.1 M_\odot \text{s}^{-1}} \right)
\times \left( \frac{\epsilon_{\text{nuc}}}{6 \text{ MeV}} \right) \text{erg s}^{-1}.
$$

To unbind the overlying layers of the progenitor with a gravitational binding energy of $\sim 0.6 \times 10^{51} \text{erg}$ would therefore require massive outflow for several 100 ms. The neutrino-driven wind after the end of accretion as well as explosive nuclear burning will provide additional energy. For a reliable determination of the final SN explosion energy, model 3Ds would need to be evolved considerably longer than in our simulation.

5. CONCLUSIONS

We showed that strangeness contributions to neutrino–nucleon scattering with an axial-vector coupling of $g_\sigma = -0.2$ are sufficient to turn a non-exploding 3D simulation of a $20 M_\odot$ model (in which $g_\sigma = 1.26$ was used for the standard isovector form factor) to a successful explosion. Strange-quark effects in the nucleon spin reduce the neutrino opacity of neutron-rich matter inside the neutrinosphere and thus directly and indirectly enhance the luminosities and mean energies of the radiated neutrinos. This leads to amplification of the neutrino-energy deposition behind the stalled shock because charged-current interactions are not affected by nucleon strangeness. Owing to a reduced neutrino-heating timescale and stronger non-radial mass motions in the gain layer, the shock is driven to runaway expansion $\sim 100$ ms after the passage of the Si/Si+O interface. The enhanced neutrino emission enabled by the strangeness effects is associated with (and partly caused by) a faster contraction of the PNS and a corresponding rise of the neutrinospheric temperatures. Strangeness corrections in neutrino–nucleon scattering are therefore similarly beneficial for shock revival as “softer” nuclear EOSs, which also lead to faster PNS contraction and stronger emission of more energetic neutrinos (Marek & Janka 2009; Janka 2012; Suwa et al. 2013).

Our results demonstrate how close previous, unsuccessful 3D core-collapse models with the PROMETHEUS-VERTEX code were to explosion. They also underline that an accurate knowledge of neutrino–nucleon interaction rates, in particular also for neutral-current scattering, is of crucial importance for assessing the viability of the neutrino-driven explosion mechanism. While strangeness contributions affect neutrino–nucleon scattering everywhere, such opacity modifications between the subnuclear regime and neutrinospheric densities are most relevant for SN shock revival. In the discussed $20 M_\odot$ simulations a modest $\sim 15\%$ diminution of neutrino–neutron scattering makes all the difference between explosion and failure. Theoretical and experimental determinations yield somewhat smaller absolute values for $g_\sigma$ than assumed in our study (Ellis & Karliner 1997; Airapetian et al. 2007; Alexakhin et al. 2007). However, in-medium effects like correlations in low-density nucleon matter may cause similar opacity reductions (C. Horowitz 2015, private communication).

Nucleon-strangeness effects should also be investigated in 3D simulations of other progenitors. Moreover, our calculations must be repeated with better than $2^\text{o}$ angular zoning to ensure that the explosion is robust and withstands higher angular resolution of the cascading of turbulent energy from the largest scales to small structures, which can be harmful for shock revival (Hanke et al. 2012; Couch 2013; Abdikamalov et al. 2014; Couch & O’Connor 2014). In any case, however, the outcome of multi-dimensional core-collapse simulations that marginally explode or fail can sensitively depend on effects on the 10% level in the neutral-current neutrino–nucleon interactions. This sensitivity needs to be taken into account in numerical implementations of these rates and might also be important for understanding partially conflicting model results published by different groups.

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