Neutrinos and Supernovae

Stephen W Bruenn[1], *Eric J. Lentz*[3,2,4], *Eric J. Lingerfelt*[5], *Anthony Mezzacappa*[2,3,5], *W. Raphael Hix*[2,3], *John N. Blondin*[6], *O. E. Bronson Messer*[7,3], *Pedro Marronetti*[1]

[1]Florida Atlantic University, 777 Glades Road, Boca Raton, FL 33431-0991, USA
[2]Physics Division, Oak Ridge National Laboratory, P. O. Box 2008, Oak Ridge, TN 37831-6354, USA
[3]Department of Physics and Astronomy, University of Tennessee, Knoxville, TN 37996-1200, USA
[4]Joint Institute for Heavy Ion Research, Oak Ridge National Laboratory, P. O. Box 2008, Oak Ridge, TN 37831-6374, USA
[5]Computer Science and Mathematics Division, Oak Ridge National Laboratory, P. O. Box 2008, Oak Ridge, TN 37831-6164, USA
[6]Department of Physics, North Carolina State University, Raleigh, NC 27695-8202, USA
[7]National Center for Computational Sciences, Oak Ridge National Laboratory, P. O. Box 2008, Oak Ridge, TN 37831-6164, USA

DOI: http://dx.doi.org/10.3204/DESY-PROC-2011-03/bruenn

We review the current status of the core-collapse supernova (CCSN) mechanism, beginning with a very brief account of CCSN types, and of the growing number of observations of their corresponding progenitors. This is followed by a brief account of current developments in the modeling of CCSNe progenitors, and of the four main supernova mechanisms under current study. We then focus on the current status of the neutrino mechanism, describing its essential features, as this mechanism seems the most promising at this time. We follow with a brief description of current efforts at the very computationally demanding challenge of modeling the neutrino mechanism, ending with a brief description of our CHIMERA code and some recent results obtained with this code.

1 Progenitors of Core-Collapse Supernovae

1.1 Observations

One of the major observational goals of core-collapse supernovae (CCSNe) research is to determine what main sequence mass range of stars will end up as CCSNe. Computer simulations of stellar evolution indicate that stars between 8 $M_\odot$ and $\sim$ 140 $M_\odot$ end their normal thermonuclear evolution with a core collapse. Below 8 $M_\odot$, a star sheds its envelope and settles non-explosively into a white dwarf. Above $\sim$ 140 $M_\odot$ a star suffers a pair-instability, triggering explosive oxygen core burning which results in the complete disruption of the star.

CCSNe are classified, in order of increased envelope stripping prior to explosion, as II-P (plateau), II-L (linear), IIn (narrow lines) IIb (transitional), Ib (no H), and Ic (no H or He) [1]. Until the advent of SN1987A, estimates of progenitor masses for any of these types had to rely on the correlations between these supernovae and regions of massive star formation (giant HII regions) in galaxies, or estimating the ejected mass by modeling light curves and ejection velocities. Because of the nearness of SN1987A (a type II pec at $\sim$ 48 kpc) it was possible for the first time to identify the progenitor (a blue supergiant of mass $\sim$ 18 $M_\odot$). In addition, the detection of a total of 19 neutrinos in the Kamioka and IMB detectors confirmed the basic CCSN scenario, namely, that a core collapse of a massive star does indeed occur, and that the bulk of the gravitational binding energy is radiated away by neutrinos. It was also possible to detect the progenitor for the next nearest explosion, SN1993 in M81 ($\sim$ 3.6 Mpc), a type II which transformed to a type Ib, the data being best fit by a star of 15 $M_\odot$ main-sequence (MS) mass and tidally stripped of most of its hydrogen envelope by prior binary mass transfer.

The detection of other supernova progenitors has been made possible by high resolution archives of galaxy images built up over the past 15 or so years, mainly by the Hubble Space Telescope. By finding an image or an upper limit to an image on a high resolution pre-explosion galaxy image at the position of the SN it is possible, by comparing the color and magnitude of the image (or upper limit) with stellar models, to estimate the MS mass (or upper limit) and stellar type of the progenitor. For the progenitors of type II-P SNe, the most common type, this technique has netted 8 identifications (all red supergiants (RSGs), as expected) for which mass estimates are possible, and 12 upper limits (see [2] for a review through early 2009). Surprisingly, within the uncertainties of stellar modeling and the possible extinction of the progenitor by dust, no progenitor of a SN II-P greater than 21 $M_\odot$ has been found, and this upper limit is probably smaller. The lower limit appears to be about 8.5 $M_\odot$. Two progenitors of SNe II-L have been identified (tentatively) as yellow supergiants (YSGs) with inferred progenitor masses of 18 - 25 $M_\odot$, and only one progenitor of a SN IIn, a massive ($\geq$ 50 $M_\odot$) luminous blue variable (LBV). Several Type IIb SN progenitors have been identified; SN1993 mentioned above, and SN 2011dh whose progenitor has been identified as YSG of initial MS mass of 18–21 $M_\odot$ [3] or 13 $\pm$ 3 $M_\odot$ [4]. Unlike the above SN types, no progenitors of SNe Ib or Ic have been identified. They are expected to arise from Wolf-Rayet stars and/or from stars whose H or H/He envelope has been tidally stripped by Roche lobe overflow to a binary companion. Their correlation with regions of massive star formation suggest that at least some arise from the former.

It thus appears from observations that single stars in the MS mass range from $\sim$ 8.5 to 16 - 21 $M_\odot$ explode as SN II-P, and perhaps a subset of stars in this mass range that are tidally stripped in binaries and more massive single WR stars explode as SNe IIb, SNe Ib or Ic. On the basis of very limited observations, more massive stars explode as SN II-P when YSG's, and very massive LBV explode as SN IIn. Unanswered are the fate of RSG more massive than $\sim$ 21 $M_\odot$, whether some stars collapse directly to black holes with little or no optical display, whether there are distinct faint and bright branches of SN from a given mass range, and so on.

1.2 Modeling

The evolutions of massive stars from the main-sequence (MS) to the onset of core collapse have been computed in 1D by a number of groups ([5] and references therein). These calculations have all shown the development of a generic stellar structure consisting of nested burning shells of successively heavier elements from the hydrogen-helium envelope on the outside to the silicon-like shell overlying an inert iron-like core at the center. One of the great uncertainties attending 1D progenitor calculations is the modeling of turbulent mixing processes—e.g., turbulent mixing associated with thermal convection and/or rotational shear. Standard procedure is to use some form of mixing length theory (MLT) in the turbulent regime with a prescription as to the location of the inner and outer boundaries, and to assume physically plausible diffusion coefficients for the transport of angular momentum due to rotational shear, magnetic torques, etc, ([6] and references therein). This introduces a number of ambiguities into the model, e.g., placement of the convective boundaries, mixing profiles, interaction with of convective plumes with the stable layers, diffusion coefficient magnitudes, etc.

It has now become possible with improved computing facilities to study the hydrodynamics of stellar interiors with two- and particularly three-dimensional simulations of space and time "windows" of the stellar interior (see [7] and references therein for modeling thermal convection, and [8] and references therein for modeling rotation). These studies have revealed shortcomings of current 1D modeling. In particular, they have shown that: (1) the traditional prescriptions for setting the locations of the inner and outer boundaries of a turbulent layer, e.g., the Schwarzschild or Ledoux criterion, should be replaced by a Richardson criterion, which is a measure of the boundary "stiffness" to the strength of the turbulence; (2) rather than convective overshooting, mixing occurs at convective boundaries and the convective region grows due to turbulent entrainment (shear instabilities and the scouring out of the stable layer by turbulent eddies); (3) wave generation at convective boundaries induces slow mixing in the stable layers; (4) shell burning, particularly O shell burning, can be violent, leading to asymmetries at the onset of core collapse; (5) smaller effective mixing lengths near the lower boundaries of a turbulent regions lead to steeper temperature gradients, affecting nuclear-reaction rates due to their stiff temperature dependence; (6) the large ratio of the size of rising high entropy plumes to low entropy downflows results in a large downward flux of kinetic energy (neglected in MLT) compensated by a correspondingly large upward enthalpy flux to maintain the requisite convective luminosity; (7) in rotating red giant envelopes, simulations do not approach the extreme cases of uniform mean radial specific angular momentum or angular velocity as has been typically assumed in 1D models.

It is clear from the above that the structure of core collapse progenitors may be expected to change in the near future as results of these multi-D computational experiments become incorporated into stellar evolutionary calculations.

2 The Core-Collapse Supernova Mechanism

Four explosion mechanisms have been the focus of CCSNe research during recent years: (1) the acoustic mechanism, (2) the MHD mechanism, (3) the hadron-quark mechanism, and (4) the neutrino-driven mechanism.

2.1 Acoustic Mechanism

The acoustic mechanism was discovered by [9] in the simulation of an 11 $M_\odot$ progenitor. They found that long after shock stagnation ($\gtrsim 0.6$ s post-bounce) turbulence and anisotropic accretion on the proto-neutron star excites and maintains vigorous g-mode oscillations which radiate intense sound waves, the energy coming from the gravitational binding energy of the accreted gas. As these sound waves propagate outward through the negative density gradient into the surroundings, they steepen into shocks and their energy and momentum are efficiently absorbed, powering up the supernova explosion. Thus, the proto-neutron star acts like a transducer converting the gravitational energy of infall into acoustic energy which propagates out and deposits energy in the surroundings, in analogy with the neutrino transport mechanism. They subsequently found [10] that the acoustic mechanism is able to explode a variety of progenitors at late times ($\gtrsim 0.6$ s). The physical reality of this mechanism is debated, as it has not been observed in simulations by other investigators although their numerical techniques, though different, are capable of capturing this phenomenon. A further note of caution is cast by a recent study by [11] that finds that the damping of the primary $\ell = 1$ g-mode mode by the parametric instability causes this mode to saturate at an energy two orders of magnitude lower than that required to power a supernova.

2.2 MHD Mechanism

Magnetic fields threading a progenitor are frozen in the gas on all relevant core-collapse time scales. On core collapse these magnetic fields will be amplified both by flux conservation during matter compression and by being wound up toroidally by the differential rotation of the core. Simulations with increasing sophistication have shown that if the iron core before collapse is threaded by very strong magnetic fields (B $\geq 10^{12}$ gauss), then this in combination with rapid rotation can produce jet-like explosions magnetically on a prompt time scale (see [12, 13, 14, 15, 16, 17]). Furthermore, it has been recognized that initially weak magnetic fields can be amplified to equipartition values exponentially by the magnetorotational instability [18, 19, 20]. Notwithstanding all this, it must be appreciated that the maximum magnetic energy that can be achieved in a differentially rotating core is the free energy, $T_{\rm free}$, of the differential rotation, i.e., the difference between the energy of the differentially rotating core and the same core uniformly rotating with the same angular momentum, and

$$T_{\rm free} \leq T_{\rm rot} = 4 \times 10^{51} \left(\frac{\kappa_{\rm I}}{0.3}\right) \left(\frac{M}{1.4 M_{\odot}}\right) \left(\frac{R}{10\ {\rm km}}\right)^{2} \left(\frac{P_{\rm rot}}{2\ {\rm ms}}\right)^{-2} \ {\rm ergs} \tag{1}$$

[20]. Thus rather small initial rotation periods, ≤ 2 ms, for newly formed neutron stars are required if enough magnetic energy is to potentially arise to power up the typical supernova. These small rotation periods are at variance with the calculated rotational periods of the magnetized cores of supernova progenitors [6], and the extrapolated periods of newly formed neutron stars (≥ 10 ms). Both of these constraints are "soft" (stellar evolutionary calculations with rotation and magnetic fields are not ab initio, and we have not yet observed a newly formed neutron star), but if they hold, then the MHD mechanism will only be relevant to a subset of core collapse supernovae. However, the observations of magnetars, long-duration gamma-ray bursts, and hints of highly collimated material in some supernova remnants suggests a subclass of events that are magnetically driven.

2.3 Hadron-Quark Mechanism

A recently investigated ([21, 22, 23]) possible supernova mechanism obtains if a hadron-quark phase transition occurs at the $(\rho, , T, Y_{\rm e})$-values sampled by the core center at and around core bounce. As in all scenarios, the core bounce launches a shock which propagates out to 100 - 200 km and stalls, becoming an accretion shock (see below). The formation of a mixed phase softens the equation of state (EOS) and induces a further collapse of the protoneutron star (PNS) at some given period after the initial core bounce when this EOS softening encompasses enough of the core. A secondary accretion shock forms when this second core collapse is halted by the formation of a pure quark phase which stiffens the EOS again. The secondary accretion shock becomes dynamic, propagates outward, and merges with the original accretion shock and the resulting shock propagates outward giving rise to a supernova. The viability of this model depends on a low critical density for the hadron-quark phase transition for proton fractions of 0.2 - 0.3 and temperatures of 10 - 30 MeV. At the present time this critical density is unknown.

2.4 Neutrino-Driven Mechanism

The neutrino-driven mechanism has a long pedigree extending back to the seminal paper of [24] and its more modern incarnation [25]. Following the collapse and bounce of the inner core of massive star at the endpoint of its normal thermonuclear evolution, the shock launched at core bounce stalls in the outer core, losing energy (and therefore post-shock pressure) to nuclear dissociation and electron neutrino losses. Within a short time (~ 50 ms) a thermodynamic profile is established in which infalling matter encountering the outward flow of neutrinos undergoes net heating between the shock and the so-called gain radius, and net cooling below, due to the different neutrino heating and cooling radial profiles. Crudely speaking, for neutrino heating to be successful in powering an explosion a fluid element must be heated sufficiently while it resides in the heating layer to reenergize the shock.

Energy deposition by neutrinos plays the primary role in the neutrino-driven mechanism, and the rate of energy deposition per nucleon, $\dot{q}$, can be written as

$$\dot{q} = \frac{X_{\rm n}}{\lambda^{a}_{\nu_{\rm e}}}\frac{L_{\nu_{\rm e}}}{4\pi r^2}\langle\epsilon^2_{\nu_{\rm e}}\rangle\frac{1}{f_{\nu_{\rm e}}} + \frac{X_{\rm p}}{\lambda^{a}_{\bar{\nu}_{\rm e}}}\frac{L_{\bar{\nu}_{\rm e}}}{4\pi r^2}\langle\epsilon^2_{\bar{\nu}_{\rm e}}\rangle\frac{1}{f_{\bar{\nu}_{\rm e}}}, \tag{2}$$

where the first and second terms express the absorption of electron neutrinos ($\nu_{\rm e}$'s) and antineutrinos ($\bar{\nu}_{\rm e}$'s), respectively. For the $\nu_{\rm e}$'s ($\bar{\nu}_{\rm e}$'s), $L_{\nu_{\rm e}}$ ($L_{\bar{\nu}_{\rm e}}$) is their luminosity, $\langle\epsilon^2_{\nu_{\rm e}}\rangle$ ($\langle\epsilon^2_{\bar{\nu}_{\rm e}}\rangle$) their mean square energy, and and $\frac{1}{f_{\nu_{\rm e}}}$ ($\frac{1}{f_{\bar{\nu}_{\rm e}}}$) their inverse flux factor (ratio of zeroth to first angular moment of the neutrino distribution), which is a measure of their anisotropy. The presence of the mean square neutrino energies and the inverse flux factors underscores the necessity of accurately calculating both the energy spectrum and the angular distribution of the neutrinos and antineutrinos.

3 Status of the Neutrino-Driven Mechanism

Simulations of core-collapse supernovae in spherical symmetry with considerable realism have been performed with Boltzmann neutrino transport, state-of-the-art neutrino interactions, and with/without general relativity [26, 27, 28, 29, 30]. Except for the smallest MS masses undergoing core collapse, these have not yielded explosions. Something was clearly missing.

An insight as to what might be the missing ingredient began to be appreciated during the 1990's, and is the essential role played by multidimensional effects. Analyses of immediate post-bounce core profiles given by computer simulations had for a long time indicated that a variety of fluid instabilities are present, driven by gradients in entropy and/or leptons. The most important of these for the neutrino-driven mechanism is the neutrino heating above the neutrinosphere. Because neutrinos heat the bottom of the heating layer most intensely, a negative entropy gradient builds up which renders the layer convectively unstable. In order for convection to grow, however, the fluid must remain in the heating layer for a critical length of time; roughly the ratio of the advective timescale to some averaged timescale of convective growth timescale must be $\gtrsim 3$ [31]. If convection can get established in the hearing layer, hot gas from the neutrino-heating region will be transported directly to the shock, while downflows simultaneously will carry cold, accreted matter to the layer of strongest neutrino heating where a part of this gas, being cold, readily absorbs more energy from the neutrinos. The loss of energy accompanying the advection of matter through the gain radius is thereby reduced and more energy stays in the heating layer. A useful criterion for shock revival and successful neutrino driven explosions is the residency time of fluid elements in the heating layer [32]. Longer residency times, by allowing fluid elements to absorb more energy, are thus favorable for explosions. The lateral motions associated with convection lengthen the heating layer residency time of some fluid elements, thereby further enhancing the effects of convection.

Another important phenomenon missing in spherical symmetry was pointed out by [33] who discovered that the stalled shock is subject to low-mode aspherical oscillations, which is now referred to as the standing accretion shock instability or 'SASI.' The cause of this instability is still being debated, but it leads to enlargements of the heating layer in certain regions and its diminution in others. Where the heating region is enlarged, the residency times of fluid elements is lengthened, and the magnitude of neutrino heating is enhanced. Where the heating region is constricted, conditions are favorable for the establishment of down-flows or return-flows for large-scale convection. The diversion of infalling material by the distorted shock to the constricted regions adds to these effects. It has further been pointed out and supported by 2D simulations with parameterized neutrino sources that the development of the SASI leads to the large asymmetries observed for SN 1987A and other supernovae, and might account for the large observed velocities of neutron stars [34, 35].

A number of groups are actively engaged in modeling CCSNe with the aim of ascertaining the viability of the neutrino mechanism. These include the Princeton-CalTech-Israeli group using the Vulcan/2D code, the Swiss-Japanese group using the Zeus+IDSA code, the Garching group with using VERTEX code, and the FAU-NCSU-Oak Ridge group using the CHIMERA code, e.g., ([36, 37, 38, 39] and references therein to earlier work). Because modeling realistic neutrino transport in CCNSe is such a computational challenge, demanding exascale resources to fully implement, say, the solution for each neutrino specie of a multi-energy, multi-angle, GR Boltzmann transport in three spatial dimensions (i.e., the solution of the time evolution of neutrinos in a six-dimensional phase space), these groups have had to resort to various approximate approaches to reduce the dimensionality of the problem. Zeus+IDSA, VERTEX, and CHIMERA use a ray-by-ray approximation whereby the three spatial dimensions accessible to a neutrino are reduced to "one-and-one-half" dimensions—namely, a one-dimensional transport problem is solved along each of an ensemble of independent radial rays spanning the solid angle of the computational domain. Lateral neutrino pressure and advection, but not transport, are computed in neutrino optically thick regions. Vulcan/2D solves the transport equation along independent "energy rays," thus ignoring non-isoenergetic scattering, specie coupling, and observer corrections, the latter omission being potentially serious as it leads, among other things, to a substantial non-conservation of energy and leptons [40].

The codes also use different approximations in solving their respective transport equations. CHIMERA and Vulcan/2D use flux-limited diffusion (the latter also has an SN option), VERTEX uses a variable Eddington tensor closure with the latter calculated by the solution of a spherically averaged, model Boltzmann equation. Zeus+IDSA uses an Isotropic Diffusion Source Approximation (IDSA) [41], whereby the neutrino distribution is broken up into a trapped component and a streaming component. VERTEX and CHIMERA use approximate GR for both the gravity and transport, CHIMERA using a spherically symmetric post-Newtonian approximation gravitational potential supplemented by a Newtonian gravity spectral Poisson solver described in [42], VERTEX having used the same until recently switching to the conformally flat approximation [43], while the other two codes use the Newtonian approximation in both gravity and transport. A full opacity set is used by both VERTEX and CHIMERA (VERTEX also includes neutrino-neutrino interactions), including non-isoenergetic scattering from nucleons, electrons, and positrons, ion-ion correlations, weak magnetism corrections, and electron capture on an ensemble of heavy nuclei using the improved data of [44]. Vulcan/2D and Zeus+IDSA omit non-isoenergetic scatterings, ion-ion correlations, and weak magnetism corrections, and treat electron captures on nuclei as described in [45]. Finally, CHIMERA independently transports four coupled neutrino species ($\nu_{\rm e}$'s, $\bar{\nu}_{\rm e}$'s, $\nu_{\mu\tau}$'s and $\bar{\nu}_{\mu\tau}$'s), VERTEX independently transports three coupled neutrino species ($\nu_{\rm e}$'s, $\bar{\nu}_{\rm e}$'s, and $\nu_{\rm x}$'s, where $\nu_{\rm x} = \{\nu_{\mu\tau}, \bar{\nu}_{\mu\tau}\}$), and Vulcan/2D and Zeus+IDSA transport two uncoupled species ($\nu_{\rm e}$'s, $\bar{\nu}_{\rm e}$'s).

Given the differences between the groups in the implementation of neutrino transport, and other differences, such as the grid geometry and resolution, the hydrodynamic scheme, equation of state used, and the progenitors chosen, it is not surprising that results from the above groups have not converged. The FNO group in has found, for example, fairly strong explosions for progenitor masses in the range of 12 - 25 $M_\odot$ [39], the G and SJ groups get, respectively, weak explosions for 11.2 and 15 $M_\odot$ progenitors [46], and a 13 $M_\odot$ progenitor [37], while the PCI group did not obtain any neutrino-driven explosions in their simulations [9, 10]. Clearly, much work needs to be done before the viability of the neutrino-driven mechanism is understood.

4 CHIMERA Code

The CHIMERA code has been extensively refined and updated since the last suite of core-collapse simulations were performed as described in [39]. Adding to the brief description of the code given above, the CHIMERA hydrodynamics consists of a Lagrangian remap implementation of the Piecewise Parabolic Method (PPM) [47]. A moving radial grid option and an adaptive mesh redistribution algorithm keeps the radial grid between the core center and the shock structured so as to maintain approximately constant $\Delta\rho/\rho$ both during collapse and the post-bounce evolutionary phases. The grid resolution for current 2D simulations is now 512 radial zones and 256 angular zones . Above 10^{11} g cm^{-3}, the equation of state (EOS) used is the Lattimer-Swesty EOS [48] with an incompressibility coefficient of 220 MeV. For matter in nuclear statistical equilibrium (NSE) below 10^{11} g cm^{-3} with proton fractions below 26/56 a modified Cooperstein EOS is used [49], and for a proton fraction above 26/56 the nuclei are treated as an ideal gas of 17 nuclear species in NSE (neutrons, protons, alpha particle nuclei from 4He to 60Zn, and 56He). In regions where nuclei are not in NSE, the same 17 nuclei are employed with the 14 alpha particle nuclei evolved by means of a nuclear reaction network.

5 2D Simulation Results

A suite of four 2D simulations are currently being computed by CHIMERA, initiated from the 12, 15, 20, and 25 $M_\odot$ core-collapse progenitors of [50]. The 15, 20, and 25 $M_\odot$ models have evolved less than 200 ms post-bounce, which is too early to ascertain whether an explosion will ensue. We will briefly describe the 12 $M_\odot$ model, which has evolved past 200 ms and seems to be exhibiting the beginnings of an explosion.

Following core bounce, which occurs 260.9 ms after the initiation of the simulation, the shock propagates rapidly out to $\sim$ 60 km in 10 ms and than pauses for a few ms. At this time a brief but violent convection sets in initially at a radius of ~ 38 km, driven by a negative lepton fraction (Y_ℓ) gradient. This gradient is established when the shock propagates through the neutrinosphere, and it occurs by the rapid electron capture and outward transport of ν_e's, which lowers the electron fraction (Y_e) and Y_ℓ below the values in the denser regions below where the ν_e's are fully trapped. The inner radius of this region of lowered Y_e is advected inward as the core compresses while the outer radius of this region advances with the shock. Consequently, in ~ 5 ms the lepton-driven convection grows to encompass the region between 16 and 60 km, the shock in the mean time having resumed its outward propagation to ~ 80 km. This convective episode only lasts only $10 - 15$ ms, but it has the effect of perturbing the shock and exciting low-order SASI modes. During the next ~ 30 ms the shock continues to propagate out and then stagnates at an average radius of ~ 150 km, all the time sloshing from side to side in what appears to be a combination of $\ell = 1$ and $\ell = 2$ modes. This sloshing motion causes entropy fluctuations in the form of arcs to be formed behind the shock. Arcs of high entropy material form behind the shock when the shock front is moving outward and the relative velocity between the shock front and the infalling material is correspondingly large; arcs of low entropy material form behind the shock when the shock front moves inward and the relative velocity between the shock front and the infalling material is correspondingly small. Comparing the mean shock radius versus time of this 2D simulation with the shock radius versus time of a corresponding 1D simulation shows that the two radii track each other closely for the first 30 or so ms. After 30 ms, the mean shock radius of the 2D simulation begins to exceed that of the 1D simulation. The shock radius of the 2D simulation also exhibits a growing difference between its maximum and minimum values, reflecting the SASI and its increasing amplitude with time. The mean entropy of the fluid as a function radius in the 2D simulation at 30 ms post-bounce tracks closely that of the 1D simulation out to about 40 km, but rises slightly above that of the 1D simulation out to the shock, with a growing spread between maximum and minimum values.

At about 50 ms post-bounce the heating layer begins to become convectively unstable, and mushroom shaped plumes of high entropy material begin to appear, separated by narrow down-flows. These plumes slowly merge, and by 90 ms post-bounce a flow pattern from polar to equatorial regions becomes established. This is caused by the prolate shape assumed by the shock at this time. Infalling material encountering the prolate shock is deflected equatorially, driving the polar to equatorial flow. By 130 ms post-bounce smaller high entropy convective plumes have merged into three large high entropy plumes, two polar and one equatorial, with down-flows on either side of the equatorial plume. The residency time for material in the high-entropy plumes, particularly the polar plumes, becomes large, while newly shocked material is directed towards the down-flows and quickly reaches the surface of the nascent neutron star. The large residency time of the material in the high-entropy plumes causes the entropy to rise there, with the result that the plumes tend to push the shock farther out. The down-flows are deflected around the nascent neutron star by the SASI induced motion of the fluid there, and encounter the neutron star surface near the polar regions causing the accretion ν_e and $\bar{\nu}_\mathrm{e}$ luminosities to peak there. The result is that at this time the ν_e and $\bar{\nu}_\mathrm{e}$ luminosity as a function of angle is generally smaller in the equatorial region than in a corresponding 1D simulation at the same time, but larger in the polar regions. By 210 ms post-bounce the shock near the zero angle polar region passes 400 km, and the shock near the near the opposite pole passes this radius 20 ms later. A runaway situation has apparently been established and the shock continues to expand in a highly prolate fashion for the rest of the simulation. The simulation has only been carried out for 240 ms so it is much too early to ascertain the explosion energy.

6 Conclusions

Observations of CCSNe during the past decade are providing information as to the masses and evolutionary states of their stellar progenitors, and the types of stellar progenitors as a function of the SN spectral type. The statistics are still poor except for the type II-P SNe, which arise from a range of initial MS masses from ~ 8.5 M$_\odot$ to an uncertain maximum mass which is likely metalicity dependent. The low end of the initial MS mass of CCSNe progenitors, in fact, appears to be ~ 8.5 M$_\odot$, as expected from theory. Computer modeling of these progenitors can be expected to become more realistic as new multi-dimensional computations of fluid instabilities, induced by thermal convection and rotation, serve to guide their implementation in the stellar models. The most extensively studied CCSN mechanism, and perhaps the most promising at this time, is the neutrino-driven mechanism. The formidable computing challenges of realistically modeling the neutrino-driven mechanism, however, has compelled groups to make approximations in the physics and the numerics implemented in their codes, particularly in their neutrino transport algorithms, and this has led to disagreements between the results of different groups. It is expected that as more powerful computing resources become available enabling more realistic modeling to be done, these disagreements will diminish. The neutrino-driven mechanism will the either become the standard CCSN model, or will be shown to be in need of modification or replacement.

7 Acknowledgments

S.W.B. and P.M. acknowledge partial support from NASA award 07-ATFP07-0011, E. J. Lentz. is supported NASA grant number NNH11AQ721 and NSF grant number OCI-0749242), A.M. and W.R.H. are supported by the Department of Energy Office of Nuclear Physics, and A.M., O.E.B.M., and E.J.Lingerfelt are supported by the Department of Energy Office of Advanced Scientific Computing Research. This research was supported in part by the National Science Foundation through TeraGrid resources provided by the National Institute for Computational Sciences under grant number TG-MCA08X010. This research used resources of the Oak Ridge Leadership Computing Facility at the Oak Ridge National Laboratory, which is supported by the Office of Science of the U.S. Department of Energy under Contract No. DE-AC05-00OR22725.

References

[1] A. V. Filippenko. *Annual Review of Astronomy and Astrophysics*, 35:309, 1997.

[2] S. J. Smartt. *Annual Review of Astronomy and Astrophysics*, 47:63, 2009.

[3] S. D. Van Dyk, W. Li, S. B. Cenko, M. M. Kasliwal, A. Horesh, E. O. Ofek, A. L. Kraus, J. M. Silverman, L Arcavi, A. V. Filippenko, A. Gal-Yam, R. M. Quimby, S. R. Kulkarni, O. Yaron, and D. Polishook. *eprint arXiv:1106.2897*, 2011.

[4] J. R. Maund, M. Fraser, M. Ergon, A. Pastorello, S. J. Smartt, J. Sollerman, S. Benetti, M.-T. Botticella, F. Bufano, I. J. Danziger, R. Kotak, L. Magill, A. W. Stephens, and S. Valenti. *Astrophysical Journa Letters*, 739:L37, 2011.

[5] S.-C. Yoon and S. E. Woosley. *Astrophysical Journal*, 725:940, 2010.

[6] A. Heger, S. E. Woosley, and H. C. Spruit. *Astroparticle Physics*, 626:350, 2005.

[7] W. D. Arnett and C. Meakin. *Astrophysical Journal*, 733:78, 2011.

[8] A. S. Brun and A. Palacios. *Astrophysical Journal*, 702:1078, 2009.

[9] A. Burrows, E. Livne, L. Dessart, C. D. Ott, and J Murphy. *Astrophysical Journal*, 640:878, 2006.

[10] A. Burrows, E. Livne, L. Dessart, C. D. Ott, and J Murphy. *Astrophysical Journal*, 655:416, 2007.

[11] N. N. Weinberg and E. Quataert. *Monthly Notices of the Royal Astronomical Society, Letters*, 387:L64, 2007.

[12] A. Burrows, L. Dessart, E. Livne, C. D. Ott, and J. Murphy. *Astrophysical Journal*, 664:416, 2007.

[13] L. Dessart, A. Burrows, E. Livne, and C. D. Ott. *Astrophysical Journal*, 669:585, 2007.

[14] H. Mikami, Y. Sato, T. Matsumoto, and T. Hanawa. *Astrophysical Journal*, 683:357, 2008.

[15] S. G. Moiseenko and S. Bisnovatyi-Kogan, G. *Astrophysics and Space Science*, 311:191, 2007.

[16] H. Sawai, K. Kotake, and S. Yamada. *Astrophysical Journal*, 672:465, 2008.

[17] T. Takiwaki, K. Kotake, and K. Sato. *Astrophysical Journal*, 691:1360, 2009.

[18] S. A. Balbus and J. F. Hawley. *Astrophysical Journal*, 376:214, 1991.

[19] S. A. Akiyama, J. C. Wheeler, D. L. Meier, and I. Lichtenstadt. *Astrophysical Journal*, 584:954, 2003.

[20] M. Shibata, Y. T. Liu, S. L. Shapiro, and B. C. Stephens. *Physical Review D*, 74:104026, 2006.

[21] I. Sagert, T. Fischer, M. Hempel, G. Pagliara, J. Schaffner-Bielich, A. Mezzacappa, F.-K. Thielemann, and M. Liebendörfer. *Physical Review Letters*, 102:article id. 081101, 2009.

[22] T. et al. Fischer. *Classical and Quantum Gravity*, 27:114102, 2010.

[23] T. Fischer, I. Sagert, M. Hempel, G. Pagliara, J. Schaffner-Bielich, and M. Liebendörfer. *Astrophysical Journal Supplement*, 194:article id. 39, 2011.

[24] S. A. Colgate and R. H. White. *Astrophysical Journal*, 143:626, 1966.

[25] H. A. Bethe and J. R. Wilson. *Astrophysical Journal*, 295:14, 1985.

[26] M. Rampp and H.-Th. Janka. *Astrophysical Journal Letters*, 539:L33, 2000.

[27] A. Mezzacappa, M. Liebendörfer, O. E. Messer, W. R. Hix, F.-K. Thielemann, and S. W. Bruenn. *Physical Review Letters*, 86:1935, 2001.

[28] T. A. Thompson, A. Burrows, and P. A. Pinto. *Astrophysical Journal*, 592:434, 2003.

[29] M. Liebendörfer, A. Mezzacappa, F.-K. Thielemann, O. E. Messer, W. R. Hix, and S. W. Bruenn. *Physical Review D*, 63:103004, 2001.

[30] M. Liebendörfer, M. Rampp, H.-Th. Janka, and A. Mezzacappa. *The Astrophysical Journal*, 620:840, 2005.

[31] T. Foglzzo, L. Scheck, and H.-Th. Janka. *The Astrophysical Journal*, 652:1436, 2006.

[32] J. W. Murphy and A. Burrows. *Astrophysical Journal*, 688:1159, 2008.

[33] J. M. Blondin, A. Mezzacappa, and C. DeMarino. *Astrophysical Journal*, 584:971, 2003.

[34] K. Kifonidis, T. Plewa, H.-Th. Janka, and E. Müller. *Astronomy and Astrophysics*, 453:661, 2006.

[35] L. Scheck, T. Plewa, H.-TH. Janka, K. Kifonidis, and E. Müller. *Astronomy and Astrophysics*, 457:963, 2007.

[36] T. D. Brandt, A. Burrows, C. D. Ott, and E. Livne. *Astrophysical Journal*, 728:article id. 8, 2011.

[37] Y. Suwa, K. Kotake, T. Takiwaki, S. C. Whitehouse, M. Liebendörfer, and K. Sato. *Publications of the Astronomical Society of Japan*, 62:L49, 2010.

[38] F. Hanke, A. Marek, B. Müller, and H.-Th. Janka. *eprint arXiv:1108.4355*, 2011.

[39] S. W. Bruenn, A. Mezzacappa, W. R. Hix, J. M. Blondin, O. E. B. Messer, C. J. Dirk, and S. Yoshida. In *Journal of Physics*, volume 180 of *Conference Series*, page 012018. IOP, 2006.

[40] E. Lentz, A. Mezzacappa, O. E. B. Messer, M. Liebendörfer, W. R. Hix, and S. W. Bruenn. *Submitted to ApJ*, 2011.

[41] M. Liebendörfer, S. C. Whitehouse, and T. Fischer. *Astrophysical Journal*, 698:1174, 2009.

[42] E. Müller and M. Steinmetz. *Computer Physics Communications*, 89:45, 1994.

[43] B. Müller, H.-Th. Janka, and H. Dimmelmeier. *Astrophysical Journal Supplement*, 189:104, 2010.

[44] K. Langanke, G. Martínez-Pinedo, J. M. Sampaio, D. J. Dean, W. R. Hix, O. E. B. Messer, A. Mezzacappa, M. Liebendörfer, H.-Th. Janka, and M. Rampp. *Physical Review Letters*, 90:241102, 2003.

[45] S. W. Bruenn. *Astrophysical Journal Supplement*, 58(4):771, 1985.

[46] A. Marek and H.-Th. Janka. *Astrophysical Journal*, 694:664, 2009.

[47] P. Colella and P. R. Woodward. *Journal of Computational Physics*, 54:174, 1984.

[48] J.M. Lattimer and D. F. Swesty. *Nuclear Physics A*, 535:331, 1991.

[49] J. Cooperstein. *Nuclear Physics A*, 438:722, 1986.

[50] S. E. Woosley and A. Heger. 442:269, 2007.