

Sensitivity of Neutrino-Driven Core-Collapse Supernova Models to the Microphysical Equation of State

A Dissertation Presented for the

Doctor of Philosophy

Degree

The University of Tennessee, Knoxville

Ryan Elliott Landfield

December 2018

© by Ryan Elliott Landfield, 2018
All Rights Reserved.

Dedication

This dissertation is dedicated to the memories of my friend Philip Bannister, my grandfather Roy Campbell, and my mother Joan Landfield, all of whom were pivotal in nurturing my love for science.

Acknowledgments

My time with the UTK/ORNL astrophysics group has been perhaps the greatest learning experience of my life, from both an academic and personal perspective. I would like to thank the members my academic family during my tenure at The University of Tennessee. Firstly, I would like to thank my major professor, Tony Mezzacappa. Tony's enthusiasm for first-principles, bedrock physics is what led me to join the UTK/ORNL astrophysics group, and I appreciate greatly his patience and commitment to my development and betterment as a scientist and researcher. He was constant in his desire to help me put my research in context with respect to the field, and has a unique talent for understanding the origins of my interest in physics. I am further grateful for his willingness to indulge my interests in theory, and in all of my questions to him on some obscure aspect of physics, from quantum mechanics to quantum gravity, Tony was able to discuss the topic as if he'd given an expert lecture on it the previous day. I can honestly say that I've grown as a physicist and researcher under his mentorship, both in terms of learning to commit to a project and in that I serve myself best by dedicating my efforts to completing my knowledge of a subject. Likewise, Raph Hix has been a great teacher and confidant. I will be forever grateful for his willingness to take the time to further my understanding with each question I approached him with, whether was it relevant to my immediate work or not. He frequently claimed that the answers to my questions were over his head, yet eventually, after realizing that ultimately all of his answers tended to be the right ones, I learned that his self-deprecation is just humility. I would like to express my appreciation to Eric Lentz for his tireless mentorship and patience, as well as for time spent discussing far-reaching ideas, both physical and philosophical. I would like

to thank Eirik Endeve for his careful instruction and attention to detail. I would also like to thank Bronson Messer, whose prodigious knowledge of esoterica, superlative candor, and enthusiasm for learning make him the apotheosis of an enlightening educator. I would like to thank Andrew Steiner for his generosity with his time, his esprit d'corps for collaboration, and his professionalism as a scientist in general. I would like to thank Fernando Schwartz for his enthusiasm and insights. I would like to thank all my committee members for taking the time to be on my committee, for reviewing my work, and for providing insights to further my knowledge. I would like to thank Austin Harris, who has been both a good friend and invaluable resource during my time in graduate school; I wish to thank him for always being available to help and for his grounded perspective. I would like to thank Chris Hartnett, whose questioning nature, clever perspective, and loyal friendship were among the best aspects of my experience as a graduate student. I would like to thank my family for supporting my education throughout my life and for sharing in my wonder and appreciation of the natural world. And last but certainly not least, my special thanks go to my wife, Susan, for her patience, generous encouragement, and unwavering support so that I could complete this work.

The work presented here was supported by the National Science Foundation Nuclear Theory Program (grant PHY-1516197).

Abstract

The aim of this dissertation is to explore the sensitivity of neutrino-driven core collapse supernova physics to the microphysical equation of state of the stellar matter. Core collapse supernovae result from the gravitational collapse and ensuing explosion of massive stars. The most fundamental question in core collapse supernova (CCSN) theory is the mechanism of shock revival leading to the observed supernova explosion. The leading scenario is that neutrinos trapped in the proto-neutron star (PNS) during its formation escape and reheat the inner shock, reviving the shockwave, which expands outward, ejecting the star's outer shells into space. Given the central importance of the delayed heating mechanism to the dynamics of CCSNe, detailed studies of the dynamics of the neutrino heating mechanism, neutrino and gravitational wave (GW) signatures, and the dependence on both neutrino-matter interaction rates and the nuclear equation of state (EOS) are needed. The first direct detection of a GW emission from a CCSN would allow for a qualitative investigation into the accuracy of existing nuclear-force models via the EOS. Hence accurate CCSN numerical models are imminently necessary to predict and interpret such GW signatures. A tandem measurement of neutrino signature modulations and GWs would provide direct information about the CCSN mechanism and further our insight into the origins of neutron star rotation and kicks. In order to enhance the computational capability of the ORNL/UTK CCSN simulations, firstly I have developed the capacity to implement different pre-processed state-of-the-art EOS tables in the ORNL/UTK multi-D SN code CHIMERA. I have developed the multi-purpose EOS/Opacity table repository WeakLib for this purpose. Secondly, I have implemented these different EOS in 1D and 2D CHIMERA to study the impact of the EOS

on the collapse, bounce, and post-bounce dynamics of the CCSNe explosion, as well as the impact on neutrino emission signals.

Table of Contents

1	Motivation	1
1.1	Outline of Dissertation	2
2	Introduction to Supernovae	3
2.1	Stellar Evolution	5
2.2	The Physics of CCSNe	7
2.2.1	Conditions Before Collapse	7
2.2.2	Collapse	7
2.2.3	Core Bounce and Shock Formation	8
2.2.4	Quiescent Accretion and Shock Stagnation	8
2.2.5	Shock Revival & Explosion Mechanism	9
3	The Nuclear Equation of State	11
3.1	Regimes of the Equation of State	12
3.2	Components of the Equation of State	14
3.3	Modeling the Strong Force Interaction	14
3.3.1	The Liquid Drop Model	16
3.3.2	Independent Particle approach	17
3.3.3	Representation of Heavy Nuclei	18
3.4	The Symmetry Energy	18
3.5	The Role of Incompressibility	20

3.6	Constraints	21
3.6.1	Theoretical Neutron Matter Calculations	22
3.6.2	Nuclear Physics Laboratory Measurements	22
3.6.3	Astrophysical Observations	23
3.7	Overview of Relevant Equations of State	25
3.7.1	Lattimer & Swesty	26
3.7.2	Baron, Cooperstein, & Kahana	26
3.7.3	STOS	27
3.7.4	DD2	27
3.7.5	Steiner, Hempel, & Fischer	28
3.7.6	FSUGold & IUFSU	28
3.7.7	NL3	29
3.8	Choices of EOS for Comparison	29
4	Numerical Methods & Physics Inputs	31
4.1	WeakLib	31
4.1.1	Implementation	32
4.1.2	External	34
4.1.3	The Tables	36
4.1.4	Table Grid	37
4.1.5	Interpolation Methods	38
4.1.6	Test Problems	40
4.1.7	Thermodynamic Consistency and Self Consistency	45
4.1.8	Extrapolation Methods	45
4.2	CHIMERA	46
4.2.1	Hydrodynamics	46
4.2.2	Neutrino Transport/Radiation Hydrodynamics	47
4.2.3	Weak Physics Inputs	47
4.2.4	Nuclear Physics Inputs	48

4.2.5	Simulation Parameters	48
5	EOS in CCSN Simulation	50
5.1	CHIMERA Results	51
5.2	Effects of EOS during Collapse	61
5.3	Effects of EOS at Bounce	61
5.4	Initial Accretion Phase	68
5.5	Onset of Convection	74
5.6	Effects of EOS by 100ms	81
5.7	Effects of EOS by 200 ms	87
5.8	Effects of EOS at 300ms	100
5.9	Effects of EOS at 500ms	104
5.10	Effects of EOS at 1 second	110
5.11	Neutrino Heating by Individual Processes	113
5.12	Summary	119
5.12.1	EOS Effects on Neutrino Luminosities & Energies	119
5.12.2	Neutron Star Characteristics	122
5.12.3	Effects of EOS on Growth of Explosion Energy	124
6	Conclusion and Outlook	128
	Bibliography	132
	Vita	143

List of Tables

3.1	Summary of EOS Nuclear Parameters	30
5.1	Comparison of EOS Effects at Bounce	63
5.2	Comparison of EOS Effects at 100 ms	83

List of Figures

3.1	Constraint of Neutron Star Mass vs. Radius, by EOS. (Steiner et al. 2013)	25
4.1	WeakLib Code Structure	33
4.2	Relative pressure error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.	41
4.3	Relative entropy error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.	41
4.4	Relative internal energy error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.	42
4.5	Relative electron chemical potential error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.	42
4.6	Relative electron chemical potential error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.	43
4.7	Relative electron chemical potential error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.	43
4.8	Inversion Errors. This graph shows the correlation between temperature and internal energy of the LS EOS when queried from the table directly versus when the temperature is inversely calculated from the internal energy.	44
5.1	1D Shock Radius vs. Time. The mean shock radii for the 1D models can be seen rising initially and then peaking near 100ms and then receding for all models.	53

5.2	2D Shock Radius vs. Time. The mean shock radii for the 2D models can be seen rising monotonically for all models.	53
5.3	Diagnostic Energy vs. Time	55
5.4	Total Energy vs. Time	55
5.5	Electron Neutrino Luminosity vs. Time.	56
5.6	Electron Antineutrino Luminosity vs. Time. The magnitude of the luminosities for electron antineutrinos in each model are slightly higher than for electron neutrinos beyond the shock burst at breakout.	57
5.7	Mu and Tau Neutrino Luminosity vs. Time. The LSBCK model exhibits a higher relative peak luminosity for the μ and τ neutrino than the other models. Also, the post-400ms period also shows relatively higher luminosity with respect to the other models.	57
5.8	Mu and Tau Antineutrino Luminosity vs. Time. Again, the magnitude of the luminosities for μ and τ antineutrinos in each model is slightly higher than for μ and τ neutrinos, and the relative luminosities for LSBCK are again higher than for the ν_e and $\bar{\nu}_e$	58
5.9	PNS Mass vs. Time.	58
5.10	Gain Accretion vs. Time.	59
5.11	PNS Radius vs. Time.	59
5.12	PNS Central Density vs. Time.	60
5.13	PNS Compactness vs. Time. The evolution of the PNS compactness, the ratio of the enclosed mass to the radius.	60
5.14	Electron Fraction vs. Enclosed Mass, at $\rho_{\text{central}} = 1 \times 10^{14}$	62
5.15	Enclosed Mass vs Radius, at $\rho_{\text{central}} = 1 \times 10^{14}$	62
5.16	Density (g cm^{-3}) vs. Mass	64
5.17	Pressure (dynes/cm^2) vs. Enclosed Mass	64
5.18	Entropy vs. Enclosed Mass	65
5.19	Temperature vs. Enclosed Mass	65

5.20	PNS Mass vs. Time, at Bounce	66
5.21	Infall velocity vs. Mass	66
5.22	Electron Fraction vs. Mass	67
5.23	Electron Neutrino Luminosity vs. Mass, at Bounce	67
5.24	Electron Neutrino Luminosity vs Time	69
5.25	Electron Neutrino Luminosity Peak vs Time, zoomed view	69
5.26	Neutrino RMS Energy vs Time. The ν_e are the solid line, the $\bar{\nu}_e$ the dash-dot line, and $\nu_{\mu,\tau}$ are the dashed line.	70
5.27	Shock vs Time	71
5.28	Proto-neutron Star Radius vs Time	71
5.29	Proto-neutron Star Mass vs Time	72
5.30	Early Electron Neutrino Luminosity vs Time	73
5.31	Early Electron Antineutrino Luminosity vs Time	73
5.32	Alpha Mass Fraction and Neutrino Heating vs. Radius, 30ms.	74
5.33	Alpha Mass Fraction and Neutrino Heating vs. Radius, 40 ms.	75
5.34	Alpha Mass Fraction and Neutrino heating vs. Radius, 50 ms.	75
5.35	Gain Radius vs Time, Post-bounce	76
5.36	SFHo Entropy, 50 ms	77
5.37	SFHx Convection, 50 ms	77
5.38	DD2 Entropy, 50 ms	77
5.39	IUFSU Entropy, 50 ms	78
5.40	FSUGold Entropy, 50 ms	78
5.41	NL3 Entropy, 50 ms	78
5.42	LSBCK Entropy, 50 ms	79
5.43	Gain Radius Mass Accretion vs. Time, 50 ms	79
5.44	Gain Volume vs. Time, 50 ms	80

5.45 Shock Radius vs Time, 1D vs 2D. The one-dimensional shock trajectories can be seen here as the dotted-dashed lines while the solid lines represent the two-dimensional model shock trajectories. The 1D model shock radii peak between 80–90 ms, and start to decline, while the 2D model shock radii are beginning to rise.	80
5.46 Entropy (log scale), 100 ms	83
5.47 Neutrino RMS Energy vs Time	84
5.48 Gain Volume vs. Time, 100 ms	84
5.49 Electron Neutrino Luminosity vs Time, 100 ms	85
5.50 Electron Antineutrino Luminosity vs Time, 100 ms	85
5.51 Heating Efficiency vs Time, 100 ms	86
5.52 $\tau_{\text{adv}}/\tau_{\text{heat}}$ vs Time	86
5.53 Shock Radius vs Time	88
5.54 Gain Mass Accretion vs Time	88
5.55 Shock Mass Accretion vs Time	89
5.56 PNS Mass vs Time	89
5.57 PNS Radius vs Time	90
5.58 PNS Central Density vs Time	90
5.59 PNS Compactness vs Time. The LSBCK model’s PNS radius reaches the highest value among the models after rebounding from bounce, hence it has the lowest compactness at 20 ms. However as its radius contracts the most quickly among the models, its compactness overtakes that of the next most compact, the SFHo and SFHx models.	91
5.60 $\tau_{\text{adv}}/\tau_{\text{heat}}$ vs Time, 200 ms	92
5.61 Heating Efficiency vs Time, 200 ms	93
5.62 Electron Neutrino Luminosity vs Time, 200 ms.	93
5.63 Electron Antineutrino Luminosity vs Time, 200 ms.	94
5.64 Mu and Tau Neutrino Luminosity vs Time, 200 ms	94

5.65 Mu and Tau Antineutrino Luminosity vs Time, 200 ms	95
5.66 Neutrino RMS Energies vs Time, 200 ms.	95
5.67 Gain Heating vs Time, 200 ms	96
5.68 $\tau_{\text{adv}}/\tau_{\text{heat}}$ vs Time, 200 ms	96
5.69 Heating Efficiency vs Time, 200 ms	97
5.70 Alpha Mass Fraction vs Radius, 200 ms	97
5.71 Neutron Mass Fraction vs Radius, 200 ms	98
5.72 Proton Mass Fraction vs Radius, 200 ms	98
5.73 Shock Radius vs Time. The two-dimensional models all undergo shock revival just after 100 ms, while the shock radii of the one dimensional models monotonically contracts for the remainder of the simulation.	99
5.74 Entropy (log scale), 200 ms	101
5.75 SFHx Accretion Velocity and Mach Number, 200 ms	102
5.76 Shock Radius vs Time, 300 ms	103
5.77 Diagnostic Energy vs Time, 300 ms	103
5.78 Timescale Ratio vs Time, 300 ms	104
5.79 Heating Efficiency vs Time, 300 ms	105
5.80 Gain Volume vs Time, 300 ms	105
5.81 PNS Mass vs Time, 300 ms	106
5.82 Entropy (log scale), 300 ms	107
5.83 Alpha Mass Fraction vs Radius, 300 ms	108
5.84 Proton Mass Fraction vs Radius, 300 ms	108
5.85 Shock Mass Accretion vs Time, 500 ms	109
5.86 PNS Mass vs Time, 500ms	110
5.87 Gain Neutrino Heating vs Time, 500ms	111
5.88 Heating Efficiency vs Time, 500ms	111
5.89 Alpha Mass Fraction vs Radius, 500ms	112
5.90 Neutron Mass Fraction vs Radius, 500ms	112

5.91 Proton Mass Fraction vs Radius, 500ms	113
5.92 Entropy (log scale), 500 ms	114
5.93 Average Shock Radius vs Time, 1 second after bounce. As many of these models are highly aspherical at late times, the average shock can be significantly less than the maximum value of the shock radius.	115
5.94 Entropy (log scale), 1 second	116
5.95 Electron Neutrino Net Heating by Absorption/Emission vs Radius, 100 ms .	117
5.96 Electron Antineutrino Net Heating by Absorption/Emission vs Radius, 100 ms	118
5.97 Proton Mass Fraction vs Radius, 100 ms	118
5.98 Neutron Mass Fraction vs Radius, 100 ms	119
5.99 Neutrino RMS Energy vs Time, Post-Bounce	121
5.100PNS Central Density vs. Time	123
5.101Proto-neutron star Mass vs Radius. The point labeled “1” is the bounce event, and the point labeled “2” is the max PNS excursion after bounce, followed by the accretion phase.	123
5.102Proto-neutron star Mass vs Central Density	124
5.103Gain Neutrino Heating vs Time, 1500 ms	125
5.104Heating Efficiency vs Time	126
5.105Total Energy vs. Time	127

Chapter 1

Motivation

The modeling of core-collapse supernovae (CCSNe) is a unique challenge given the breadth of the physics involved: weak force interactions, radiation transport, nuclear microphysics, novel hydrodynamics, general relativity, etc. all need to be taken into account to provide an accurate model of the event. Furthermore, all of these constituent aspects of the model are interconnected, and all are areas of diverse research in their own right, so as improvements and revisions are incorporated into each component aspect of the physics, the others are potentially impacted as well. After debate about the fundamental mechanism driving the core-collapse explosion, it is likely that the explosion is energized by the gravitational binding energy of the compact remnant, in the form of neutrinos at first trapped inside but gradually radiated from the proto-neutron star. The formation and evolution profile of the proto-neutron star is highly dependent upon the equation of state, and hence the success or failure of accurate modeling of the explosion mechanism, which hinges on accurately describing the neutrino transport, in turn is highly sensitive to changes in the EOS quantities. Recent developments in the fields of nuclear theory as well as astronomical observation of neutron stars have allowed the production of improved modern equations of state, so it is our aim to implement these new tabular equations of state into our supernova simulations for the purposes of examining the effects of new EOS's on collapse, the formation and evolution

of the PNS, and ultimately on the dynamics of explosion, so as to better constrain and characterize our understanding of the core-collapse supernova problem in general.

1.1 Outline of Dissertation

In chapter 2, I will discuss the dynamics of CCSNe, as well as historical observations of CCSNe. In chapter 3, I will discuss the microphysics involved in our simulations, as well as the variety of different EOS used in our test simulations. In chapter 4, I will discuss the EOS/Opacity framework, WeakLib, I have developed to create, test, and implement these EOS for use and comparison in CHIMERA. Additionally, I will then discuss our CCSNe simulation code, CHIMERA, the physics and numerical methods contained therein, and our test methodology. In chapter 5, I will present my results, and in chapter 6 I will state my conclusions.

Chapter 2

Introduction to Supernovae

The collapse and death of a massive star is easily one of the most complex astrophysical events in terms of the diversity of forces and regimes that contribute to the dynamics of the event. As a star evolves over its life cycle, its mass ultimately determines its destiny, in terms of both its lifetime as well as its ultimate fate. A supernova can outshine its entire host galaxy, and ejects the majority of its stellar material into interstellar space. This is seen as an expanding cloud led by a shock wave that impacts any interstellar material it encounters, forming a nebula known as a supernova remnant. Supernovae are one of the main drivers of chemical evolution of the universe. With the exception of hydrogen, helium, and lithium produced during the Big Bang, all heavier elements are produced either in stars or during the nucleosynthesis that occurs during the supernova explosion. Furthermore, understanding of the nature of compact objects like neutron stars and black holes must involve understanding the evolutionary trajectory on which such objects evolved into their final states.

As fascination with the night sky is one of the oldest human interests, it is likely that supernovae have been observed by human eyes long ago, before early civilizations began to perform methodical and codified observations of the celestial sphere. There are inconclusive indications that peoples indigenous to the southern hemisphere witnessed and subsequently drew cave art motivated by the observation of the supernova that produced the Vela remnant

between 10,000 and 20,000 years ago. Court-appointed astronomers in ancient China recorded the appearance and persistence of “guest stars” in 185 AD in the constellation Circinus and again in 1054 AD in the constellation Taurus, the brightness and duration of which match what would be expected from galactic core-collapse supernovae (Clark and Stephenson 1982). These observations have since been confirmed (via infrared observations by Switzer and WISE, and X-ray observations via Chandra) to coincide with the remnants of supernovae 9100 and 6500 light years distant. The 1054 event was likewise recorded by Japanese and Arab astronomers, and possibly by European and Native American observers as well. This latter event gave rise to the supernova remnant we now know as the Crab Nebula. In 1572 AD, a supernova was observed by many early skywatchers, among them the Danish astronomer Tycho Brahe, and has come to be known as Tycho’s Supernova. Less than a century later in 1604 AD, another supernova was observed widely, and was recorded and studied by the younger colleague of Brahe, Johannes Kepler, and became to be known as Kepler’s Supernova. After the 1604 observation event, the next observation of a semi-galactic supernovae was the 1987 AD observation of a supernova in the Large Magellanic Cloud, now known as SN 1987A. It was observed both optically and by the detection of neutrinos, the latter of which occurred several hours before the photonic signal arrived.

Approximately 2 to 3 galactic supernovae are expected per century, though many are obscured by interstellar gas and dust, with hundreds to thousands of extragalactic supernovae observed yearly (Clark and Stephenson 1982). Observationally, supernovae are classified based on their spectra, in particular the presence or absence of hydrogen in the spectra. Supernovae with no hydrogen lines in their spectra are classified as Type I, while those with hydrogen are classified as Type II. Further grouping focused on the presence or absence of helium lines in the spectra. As more supernovae were observed, it became apparent that different types of supernovae were occurring due to fundamentally different physical scenarios, namely via either the thermonuclear detonation of a white dwarf star accreting from a binary companion, or via a second mechanism known as “core-collapse” supernova explosion. Understanding these scenarios begins by understanding the stellar life cycle.

2.1 Stellar Evolution

All stars by definition fuse hydrogen into helium. When the supply of hydrogen fuel is exhausted, the star can evolve in a number of ways, depending on the star's particular mass and chemical composition. The least massive stars fuse hydrogen to helium, but have insufficient mass to compress the resulting helium fuel to the point of fusion. The cores of stars of intermediate mass, between approximately 0.5 and 8 solar masses ($M_{\odot}$) will begin to compress until the temperature and density are sufficient to fuse helium into unstable beryllium; when helium is being fused faster than beryllium can decay, the beryllium and helium fuse into carbon and oxygen via the triple- α process. Stars in this mass range, depending on their mass and point in evolution, can fuse hydrogen in a layer outside the helium core, while fusion of helium into carbon and oxygen can eventually lead to a degenerate carbon-oxygen core. Such stars end their lives as white dwarfs once their thermonuclear evolution has ceased, and they emit radiation from dissipative cooling rather than fusion. When a white dwarf star occurs in a multi-star system, it can accrete mass from a companion star such that it approaches the what is known as the Chandrasekhar mass limit. This is the point at which electron degeneracy pressure (the emergent pressure against compression from electrons of the same spin occupying the same quantum state) in the star's core can no longer prevent gravitational collapse. Approaching this point, compression raises the carbon core density sufficiently to ignite runaway thermonuclear burning and create what is known as a thermonuclear supernova, responsible for the observational Type Ia supernova.

The post-main sequence evolution of massive stars culminates in their inevitable death via gravitational collapse and cataclysmic supernovae explosion. Massive stars having masses greater than 8–10 $M_{\odot}$ begin to form concentric shells of successively heavier elements (heavier than carbon and oxygen) in the core as the critical density and temperature required for fusion of each successive element is achieved. Carbon burns to form neon and sodium, then the neon ignites burning to oxygen, oxygen then burns to form silicon, until lastly, silicon burns to form iron. This eventually leads to the formation of a core of iron nuclei at the end

of the star’s thermonuclear evolution, which have the highest binding energy per nucleon. As the star’s stability is determined by exothermic thermonuclear burning, maintaining the pressure to keep it stable against gravitational collapse, once the iron core begins to form, nuclear burning becomes endothermic, and hence thermal pressure ceases to support the inner part of the star against gravitational collapse. For a time, electron degeneracy pressure is sufficient to keep the iron core stable, but this phase is on the order of one day as continued burning above the core is continuing the production of iron neighboring nuclei, hence adding to the gravitational force. At this point in such a star’s life, it is what we refer to as a core-collapse supernova progenitor. Depending on its composition and mass, the particular path of explosion and elemental evolution it takes may differ, but all such stars are bound at this point to begin core collapse and potentially explode.

The discovery of the neutron in 1932 first gave insight into the structure of nuclear matter, namely that the nucleus was composed of protons and neutrons. Two years after this discovery, it was suggested by Fritz Zwicky and Walter Baade that supernovae might in fact be the transition of a normal star to a novel concept, a “neutron star”, i.e. a star composed entirely of nuclear matter ([Baade and Zwicky 1934](#)). Such an object was first discovered as a pulsar in 1967 by Jocelyn Bell Burnell and Antony Hewish, and since then, thousands of others have been detected ([Hewish et al. 1968](#)). There remains a degree of uncertainty in the exact nature of the transition of a massive star to a neutron star due to theoretical and experimental uncertainties in our knowledge of the weak and strong nuclear forces, and it remains a challenge to understand exactly why some massive stars undergoing core collapse are sufficiently energetic to become supernovae while others are not. In the next few sections, I will describe the transition of a massive star to a compact remnant and explosion via the core-collapse supernova mechanism.

2.2 The Physics of CCSNe

2.2.1 Conditions Before Collapse

As stars of greater than 8–10 $M_{\odot}$ approach the end of their lives (the burning of silicon in their cores), their mantle consists of concentric layers of nuclear fusion burning composed of the remnants of the previous burning phases of its life cycle, namely hydrogen, helium, carbon, neon, and oxygen, all sustained by thermal and radiation pressure. Fusion of any element heavier than iron ceases to be exothermic and becomes endothermic. Hence when the star accumulates an iron core from the fusion burning of silicon, the core, stabilized by Fermi pressure among electrons, increases in mass until the Chandrasekhar mass limit of about 1.2 $M_{\odot}$, at which point electron degeneracy pressure becomes insufficient alone to sustain the core's stability against gravitational collapse.

2.2.2 Collapse

The silicon burning layer continues to add mass to the iron core and pushes its mass across the Chandrasekhar mass limit. As it approaches this limit, the gravitationally driven collapse of the iron core begins, with collapse of different characters occurring in the inner and outer core. The only energetic pathways available to degenerate electrons are capture on iron-peak nuclei and nucleons, which in turn lowers electron fraction and density. This serves to further reduce the local electron degeneracy pressure within the core, accelerating the collapse and favoring the production of increasingly neutron rich material in the local distribution of nuclei. This effect is strongest in the inner core, which collapses homologously, i.e. $v/R = \text{constant}$, with the outer core and all material above it in free fall, i.e. $v/\sqrt{R} = \text{constant}$ (Bethe and Wilson 1985). The inner radius of the outer core is determined by the capacity for sonic connection in the collapse timescale. Collapse proceeds, with the core losing energy to electron capture, beta decay, and photodisintegration of heavy nuclei, further accelerating collapse and releasing the gravitational binding energy of the core via vast quantities of

neutrinos (ultimately $\sim 1 \times 10^{53}$ ergs). This free escape of neutrinos continues until the core approaches $\rho \approx 1 \times 10^{12} \text{ g cm}^{-3}$ at which point the diffusion timescale of the neutrinos becomes longer than the timescale of collapse, and the neutrinos effectively become trapped in the core.

2.2.3 Core Bounce and Shock Formation

Collapse continues until nuclear densities are approached ($\sim 2.3 \times 10^{14} \text{ g cm}^{-3}$), at which spacing the strong force interactions among nucleons become repulsive, which causes the equation of state to stiffen. At this point, collapse of the inner core slows down and halts, as the inner core is overcompressed. When the inner core is no longer compressible by gravitational collapse, the remaining infalling outer core material “bounces” outward off the rebounding ultra-dense inner core. This creates a shockwave at the boundary between the homologously collapsing inner core and the supersonically collapsing outer core that expands outward. It was initially thought that this “prompt” shock was capable of ejecting the star’s outer shells and part of the core into space; such an explosion mechanism is referred to as the “prompt” mechanism. (Colgate and White 1966). When the shock reaches regions with a density of $\rho \approx 1 \times 10^{12} \text{ g cm}^{-3}$, the diffusion speed of the neutrinos overtakes the speed of the shock and a huge burst of electron neutrinos escapes through the comparatively neutrino-transparent infalling material. This robs the inner regions of significant energy.

2.2.4 Quiescent Accretion and Shock Stagnation

After the neutrino burst at shock breakout, the shock continues to expand outward while infalling matter is accreting through the shock and onto the PNS below. As the shock moves through and heats the mantle of the outer core, high temperature behind the shock causes the photodisintegration of accreting heavy nuclei (as they pass through the shock) into free protons, neutrons, and alpha particles, creating new opportunities for electron capture on

protons. reducing the energy of the shock and further transferring energy out of the matter behind the shock. The neutrino burst at shock-breakout further favors neutron-rich material and hence the deleptonization of the material behind the shock. These two sources of energy loss are sufficient to lead to the stalling of the shock against the ram pressure of the infalling material (the rest of the star), at which point the shock becomes a standing accretion shock at a radius of approximately 150 to 200 km (Bruenn et al. 2013, 2016; Janka 2012).

2.2.5 Shock Revival & Explosion Mechanism

Clearly, the shock must be revived, at least in many stars, for supernovae to be observed. The leading candidate scenario for shock revival and explosion is that neutrinos trapped in the proto-neutron star (PNS) during its formation escape and reheat the inner shocked matter, reviving the shockwave (Bethe and Wilson 1985). In this scenario, the post-core-bounce nascent neutron star contains a large number of degenerate neutrinos and antineutrinos, trapped because their mean free paths in the dense nuclear matter are significantly shorter than the radius of the PNS. Over the course of several seconds, the trapped neutrinos and antineutrinos diffuse out to the neutrinosphere, the surface of last scattering. In the explosion scenario referred to as the “delayed”, or “neutrino-reheating” mechanism (Bethe and Wilson 1985), the stalled shock wave can be thermally revived by the neutrinos streaming off the neutrinosphere. As the proto-neutron star cools via the emission of neutrinos, the neutrinos simultaneously heat the matter that has fallen through the shock. The radius at which the heating becomes longer than the cooling is known as the gain radius or gain surface in multi-D. The stellar fluid between the gain radius and the shock is heated by the neutrinos via predominantly neutrino captures on free nucleons. This heating initiates convection of the material in this region, known as the “hot bubble”. This delayed energy deposition in the stellar fluid causes the pressure to rise behind the shock, with heated up-flows rising and eventually reaching the shock. This convective engine can become sufficient to re-invigorate the shock and revive its expansion, leading to the supernovae explosion. The explosion

releases on the order of 10^{53} erg in neutrinos, originating from the gravitational binding energy lost by the collapsing iron core as it transitions to a neutron star. The neutrino energy may be what revives and sustains the shock, but the contribution of this energy to the shock is sensitive to both the physical distribution of neutrino trajectories as they interact with material behind the shock, as well as the distribution of neutrino energies and fluxes (Janka and Müller 1996; Messer et al. 1998; Janka 2001; Mezzacappa et al. 1998a). As the convective engine leads to expansion, $\approx 10^{51}$ erg of energy must be transferred to the ejected material.

Early two dimensional models introduced the capacity for convective overturn in the hot bubble, and allowed for new ways for differences in the EOS to potentially affect the reheating mechanism, with many models leading to explosion (Herant et al. 1994; Burrows et al. 1995), though the inclusion of more realistic neutrino transport led to explosion not always being achieved (Mezzacappa et al. 1998b).

The computational discovery of a new instability of the accretion shock to non-radial perturbation was shown capable of causing large scale bipolar sloshing of the shock (Blondin et al. 2003). The realistic modeling of this feature in the shockwave, known as the Stationary Accretion Shock Instability (SASI), has been shown to be essential in correctly characterizing the shock behavior and explosion dynamics (Blondin et al. 2003; Janka et al. 2007; Burrows et al. 2006; Ohnishi et al. 2006). The interplay between compression, convection, and rotational dynamics that occurs during core collapse has been suggested to be capable of jet-driven CCSN explosions (Symbalisty 1984; Akiyama et al. 2003; Burrows et al. 2007). It appears likely that the multi-dimensional hydrodynamical convective and SASI processes are crucial ingredients of the mechanism that causes the explosion of CCSN. The SASI moves the shock outward, and hence increases time matter stays in the convective gain region (Blondin and Mezzacappa 2006; Marek and Janka 2009). Through this, the SASI strengthens the deposition of energy by neutrinos to the hot bubble and ultimately assists the initiation and persistence of the explosion.

Chapter 3

The Nuclear Equation of State

In order to describe the evolution of stellar collapse, we must consider our choice of nuclear equation of state. The equation of state is the thermodynamic relation between the state variables of the stellar matter, specifically how the density, temperature and electron fraction affect, firstly, the thermodynamic quantities such as pressure, entropy, and internal energy density, and, secondly, through the nuclear composition (and the binding energies of the nuclei therein), chemical potentials, and the neutrino interactions.

Because the nuclear matter in the stellar fluid is composed of a variety of different component particles, an ensemble of nuclei and free nucleons, when the specific hybrid composition of the matter is taken into account, the equation of state provides the nuclear physics input to the supernova simulation across that ensemble. Additionally, the contribution of the additional components of the stellar medium, photons and leptons, must be taken into account. Photons are treated as an ideal Bose gas, as they are trapped in the stellar fluid. Their radiation pressure contribution is handled in this manner. Electrons and positrons are treated as a degenerate Fermi gas.

The interactions between neutrinos and the stellar fluid are governed by the weak nuclear force. The interactions between the protons, neutrons, and nuclei are governed by the strong nuclear force, the description of which, Quantum Chromodynamics, contains many

phenomenological constants. From our incomplete knowledge of the interactions comes much of the uncertainty in our description of dense matter; also contributing are the limitations in the methods for treating the many-body problem for particles interacting via the strong force. Our understanding of the strong force is still improving with better experimental nuclear data, which in turn is used to modify our current understanding of nucleon–nucleon interactions. Commensurately, different treatments and models, phenomenological or otherwise, have been implemented in astrophysical equations of state, in order to bridge the gap between our understanding of the nuclear force and its integration in the equation of state in the stellar fluid. Each equation of state I discuss will handle certain aspects of strong force interaction and many-body treatment differently, leading to the differences that will be seen in the simulation outcomes. First, these differences will be discussed in detail in the background section, as well as the required ranges of the thermodynamic phase space for the simulations. Second, I will discuss the different treatments and components that make up the equation of state for use in astrophysical simulations, and then discuss the historically relevant equations of state as well as newer equations of state used in this comparison. Third, I will discuss how we can use experimental results from nuclear physics as well as astrophysical neutron star observations to investigate which equations of state can be ruled out, and finally I will discuss the known impacts of the nuclear equation of state on the supernova event.

3.1 Regimes of the Equation of State

The equation of state must be capable of describing the dynamics of the fluid in several fundamentally different physical regimes. In stellar matter conditions below nuclear densities, the fluid is a mixture of photons, leptons, nucleons, and nuclei, the properties of each of which contributes uniquely to the dynamic evolution of the fluid. In supernova simulations, ensembles of nuclei and the reactions between them are handled by the nuclear reaction network. In high temperature and high density conditions found in the interior of CCSNe,

the nuclear reactions, governed by 3 fundamental forces, occur in these networks on widely different timescales. When strong and electromagnetic interactions of nuclei come into equilibrium at high temperatures the strong and electromagnetic reactions are typically much faster than the weak reactions, and the nuclear abundances are no longer sensitive to individual reaction rates except for the weak reactions that interconvert neutrons and protons. At these temperatures, the nuclear abundances rather depend only on the density, temperature, and the neutron fraction (which records the activity of weak reactions), and abundances in this regime can be determined by the Saha equation:

$$X_i = \left(\frac{g_i m_i^{5/2}}{\rho} \right) \left(\frac{kT}{h} \right)^{3/2} \exp \left[\frac{Z_i \mu_p^k + N_i \mu_n^k + B_i}{kT} \right] \quad (3.1)$$

where the mass fraction, X_i , for species i (${}^A_i Z_i$), on the left hand side is determined by Equation 3.1, where k is Boltzmann's constant, m_i is the mass of the particle, h is Planck's constant, N_i the neutron number, Z_i the proton number, and B_i is the binding energy. Fluid in this state is described as being in Nuclear Statistical Equilibrium (NSE).

The equation of state must describe conditions below, across, and above the temperature and density where NSE can be applied, and also in conditions approaching and above nuclear densities ($\rho \geq 2.3 \times 10^{14} \text{ g cm}^{-3}$) in the proto-neutron star. The transition region in density and temperature where free nucleons, heavy nuclei, and heavy nuclear clusters exist simultaneously is known as inhomogeneous nuclear matter. As the density becomes higher, heavy nuclei continue to exist to higher temperatures, until nuclear density is approached. The transition from the inhomogeneous regime to homogeneous nuclear matter regime is characterized by the interplay of short-range nuclear interactions and longer-ranged electromagnetic interactions becoming of similar magnitude, leading to the formation of complex nuclear structures known as “nuclear pasta”, over a narrow range of density leading to homogenous matter. While equations of state that deal with these structures are necessary for the accurate description of the interior of cold neutron stars, they are less important for the description of early-time behavior of CCSNe, where the temperature is higher and the

density is lower than in cold neutron stars. There are a great number of extant equations of state which can describe the conditions of cold neutron stars, but until recently very few that could span the regimes of temperatures, densities, and composition required for a supernova simulation. Temperatures in a CCSNe simulation range from 0.1 GK in the outer portion of the star up to 1.8×10^3 GK inside the edge of the PNS. Densities likewise range from 1×10^3 g cm⁻³ or less, up to as high as 1×10^{15} g cm⁻³. The composition, specifically the electron fraction, needs to cover the phase space from an electron fraction as low as 0.01 for nearly pure neutron matter up to 0.7 for extremely proton-rich matter which can occur in hot outflows.

3.2 Components of the Equation of State

Astrophysical equations of state vary in the exact nature of how they describe the elements of the stellar fluid and the interactions between them. Firstly, they vary in how they represent the strong force interaction between nuclei and nucleons. Additionally, different equations of state make different assumptions about the structure of nuclei in different regimes of density, temperature, and electron fraction; many equations of state are based on nucleons interactions entirely. Additionally, they vary in the treatment of nuclei and nucleons in the transition from nuclear density to inhomogeneous matter. In order to illustrate these differences, as I describe each equation of state, I will discuss how each of these aspects are deployed in the particular equation of state.

3.3 Modeling the Strong Force Interaction

Describing the strong interaction accurately from low to high density and from symmetric to heavily isospin $(n_u - n_d)/2$, where n_u and n_d are the number of up and down quarks, respectively, asymmetric matter can make the description of the most relevant physics a

complicated challenge. Ideally, such a description would arise directly from the sum of all Quantum Chromo-Dynamical (QCD) descriptions of the strong force interaction between gluons and component quarks of all nucleons and nuclei in the fluid, and QCD itself would contain no phenomenological underpinning. However, the latter is not the case and likewise the former is computationally impossible, and the current state of QCD theory doesn't provide for complete description of such interactions at the temperatures, densities, and compositions encountered in a CCSNe. Because of this, it is necessary to employ models of nuclear structure, nuclear interaction, and many-body interaction in order to reasonably describe the stellar fluid and the interactions within it. In order to minimize the uncertainty in the current knowledge of the nuclear interactions, these models are informed and constrained to the greatest degree possible by laboratory experiments and astrophysical observations in order to refine and differentiate the models and nuclear parameters. The models for the description of nuclei, the description of the nuclear interaction, and the choice of many-body method aren't necessarily independent. Here I will describe several different models and methods used to implement the strong force interaction in astrophysical equations of state, and after that the main treatments for the many-body problem for strongly interacting particles.

The approaches to model the nuclear interaction and the many-body problem can be for the most part divided into two main categories: “Ab-initio many-body methods”, and “phenomenological methods”.

1. Ab-initio methods start from the best description of inter-nucleon forces possible, but are limited in the total number of interactions they can describe without becoming intractable. Two-nucleon potential calculations are sufficient to describe the bulk of the nuclear few-nucleon observables up to intermediate energies (such as nucleon-nucleon scattering), but for three-nucleon systems (such as the binding energy of the triton), three-nucleon calculations provide a more accurate prediction of observables compared to the sum of the two-nucleon interactions. Methods used include the Hartree-Fock method, including

the relativistic (Dirac-)Bruekner-Hartree-Fock, using two-body forces based on nucleon-nucleon meson-exchange models, as well as variational and Monte Carlo methods (Oertel et al. 2017).

2. Phenomenological models and methods use effective nucleon-nucleon interactions in the place of realistic first principles methods, and are often commensurately simpler. Most of the equations of state tested in this dissertation use the phenomenological method known as the mean-field approximation, specifically the subgroup of such approximations known as relativistic mean field (RMF) approximations; non-relativistic implementations can fail in characterizing dense matter as they can create superluminal conditions. These models usually depend on parameters fitted from nuclear scattering data and the properties of nuclei and nuclear matter. Examples of such models would be the Skyrme interaction model for non-relativistic cases and, in the RMF case, relativistic meson-exchange models. Many modern effective models are expressed in terms of energy density functionals (EDF). Nuclear EDFs are commonly derived from an effective Hamiltonian solved via first order perturbation theory using Hartree-Fock methods, but are not directly connected to any nucleon-nucleon interaction. Rather, the descriptions of the microphysics are embedded in the parameters of the density-dependent forces (Oertel et al. 2017). These models will be discussed more explicitly in the coming sections.

3.3.1 The Liquid Drop Model

One of the first proposed nuclear structure models was the “liquid-drop model”, proposed by George Gamow in 1931, in which the nucleus is described as a fluid of protons and neutrons. In this model, the binding energy is determined by the balance of the electromagnetic repulsion between the protons and the strong force interaction between the nucleons, with additional contributions from the surface energy, spin pairing, and the proton/neutron asymmetry of the fluid. The binding energy of a nucleus is defined by the semi-empirical

mass formula as follows:

$$E_B = a_v A - a_s A^{2/3} - a_c \frac{Z^2}{A^{1/3}} - a_A \frac{(N - Z)^2}{A} - \delta(A, Z) \quad (3.2)$$

where Z is the number of protons and N is the number of neutrons, giving a total number of nucleons $Z + N = A$ (Weizsäcker 1935). The first term on the right-hand side of the equation is the volume term (the nuclear force interaction), followed by the surface energy term (a geometric correction to the strong force interaction), the Coulomb repulsion term, the Pauli-exclusion term (relating to proton/neutron asymmetry), and lastly the spin pairing term. The coefficients are determined in practice by the latest experimental nuclear physics data, though they can be estimated based on theory alone. In modern use this model is known as the finite range droplet model (FRDM), and reproduces experimental values for binding energies quite well, to the point that it is used often for estimation of binding energies for nuclei that have yet to be experimentally measured (Möller et al. 1995, 1997).

3.3.2 Independent Particle approach

Another approach to describe the behavior of the nucleons in the nuclear matter is the Independent Particle approach (see Oertel et al. (2017)). In this method, instead of macroscopically modeling the nucleus directly, each nucleon is treated as moving inside a potential independently from other nucleons, and is kept bound via this potential. RMF models handle many-body interactions in this way not by summing up the mutual two-body interactions of the constituent particles, but instead to describe the interaction of one particle with the remaining ones by an average potential created by the others. This is implemented in practice using the Walecka-type model, starting by defining a Lagrangian to describe the interaction terms of meson exchange, and then solving the resulting coupled set of equations. Such models are constrained by a range of parameters dictated by the experimentally verified properties of known nuclei and nuclear matter. Different parameter sets based on increasingly modern experimental nuclear data, as well as the inclusion of more

complicated interaction terms in the Lagrangian, provide the different RMF models in the equations of state included in the simulations discussed.

3.3.3 Representation of Heavy Nuclei

An additional degree of freedom in the choice of EOS is the representation of nuclei heavier than the alpha particle. All EOS's discussed here include protons, neutrons, and alpha particles as independent species, but older EOSs use what is known as the “Single Nucleus Approximation”, in which the distribution of all nuclei heavier than the alpha particle are represented as a single heavy nucleus (Lattimer et al. 1985; Lattimer and Swesty 1991; Shen et al. 1998). For the most part, this approximation reproduces the thermodynamic properties of the stellar matter quite accurately (Burrows and Lattimer 1985). However, the SNA causes an underestimation of the potential for explosion in CCSNe due to underrepresentation of the full ensemble of nuclei, which would provide increased chances for nuclear electron capture (Hix et al. 2003; Langanke et al. 2003). The EOS's included in this study other than LS employ a different method for representing the interactions with heavy nuclei. As opposed to the SNA, a fitted statistical model of nucleon interactions, as well as for nuclei and non-uniform nuclear matter, is used (Hempel and Schaffner-Bielich 2010). As opposed to the equations of state that use the SNA, these equations of state reflect the use of detailed nuclear compositions, with as many as 10000 individual species of nuclei for low-density matter. Electron capture rates have proven sensitive to this more detailed composition, as differences in nuclear structure can change electron or neutrino capture rates by orders of magnitude between nuclear neighbors.

3.4 The Symmetry Energy

One of the primary quantities that affects and characterizes the properties of nuclear matter is known as the symmetry energy, defined as the difference between the energy per baryon

of pure neutron matter and that of symmetric nuclear matter with equal numbers of neutrons and protons. Knowledge of the symmetry energy, which is density dependent, is one of the best connections between theoretical nuclear matter calculation and terrestrial experimentation of nuclear matter. Neutron-rich, highly isospin asymmetric conditions require the use of the symmetry energy to accurately describe the behavior of the stellar fluid in that regime, which is frequently encountered in supernova simulations, as well as the growth and development of the proto-neutron star. In order to calculate the symmetry energy, we expand the energy per nucleon of nuclear matter

$$E(\rho, \alpha) = E_0(\rho) + E_{\text{sym}}\alpha^2 + \mathcal{O}(\alpha^4) \quad (3.3)$$

in terms of the asymmetry parameter $\alpha = \frac{\rho_p}{\rho}$, where ρ_p is the proton density. Looking in particular at the energy per nucleon of symmetric matter

$$E_0(\rho) = m_{\text{nuc}} - B_{\text{sat}} + \frac{1}{2}Kx^2 + \frac{1}{6}Qx^3 + \dots \quad (3.4)$$

and the symmetry energy

$$E_{\text{sym}}(\rho) = J + Lx + \frac{1}{2}K_{\text{sym}}x^2 + \dots \quad (3.5)$$

where B_{sat} is the binding energy at nuclear saturation density, K is the incompressibility, Q is the skewness parameter, and L is the slope of the symmetry energy with respect to density. K_{sym} is the symmetry incompressibility, and $x = \frac{\rho - \rho_{\text{sat}}}{3\rho_{\text{sat}}}$ is the deviation of the baryon density ρ from the saturation density ρ_{sat} (see, e.g. [Oertel et al. \(2017\)](#)). The density dependence of the symmetry energy term is pivotal in connecting experimental laboratory data and the astrophysical equation of state in CCSNe. It affects many facets of CCSNe evolution, such as the neutronization of the stellar matter, as well as several aspects of PNS formation, such as the evolution of the PNS radius and the thickness of the PNS crust ([Steiner et al. 2013](#)). The density dependence of the symmetry energy is one of the key quantities that determines the relative rate of neutrino cooling in the proto-neutron star ([Lattimer](#)

and Prakash 2006). Additionally, it is involved with the favorability of the conditions for r-process nucleosynthesis in CCSNe (Lattimer and Lim 2013).

3.5 The Role of Incompressibility

The characterization of incompressibility in a given astrophysical equation of state is often described in terms of relative “stiffness” or “softness” relating the collapsed inner core to a spring. The analogy continues that for a stiff EOS, a small increase in density corresponds to a large increase in pressure relative to a softer EOS, as a stiff spring provides more force for the same displacement than a softer spring. The nuclear situation is complicated, however, as the nature of incompressibility changes between the thermodynamic region at and below nuclear saturation density and the thermodynamic region of suprasaturation nuclear density. Up to and including nuclear saturation density the symmetry energy J and its density dependence L affect the incompressibility of the nuclear matter. At supra-saturation densities the skewness parameter Q plays an increased role. The incompressibility is related to the density curvature of the energy density

$$K \sim \frac{\partial P}{\partial \rho} \sim \rho \frac{\partial^2 E}{\partial \rho^2}, \quad (3.6)$$

where ρ is the density of nucleons or baryons. The stiffness is often discussed in terms of the adiabatic index Γ_s where

$$\Gamma_s = \frac{\partial \ln P}{\partial \ln \rho_s}, \quad (3.7)$$

where Γ_s describes the stiffness at constant entropy. These quantities provide a useful characterization of the impact of the EOS on the PNS and its subsequent evolution after collapse and bounce, via the onset and behavior of convection, as a more compressible EOS gives a smaller PNS, as well as more dense stellar mass in the gain region. This additionally greatly affects the overall neutrino luminosity of the PNS and the location of the neutrinospheres. Grossly speaking, the stiffer the EOS, the larger PNS radius for the same mass, and likewise the lower central density for the same mass, as a stiff EOS is harder

to compress and offers more support against gravity. The incompressibility plays a key role in connecting calculation and experiment with astrophysical observation, specifically the extremal properties of the PNS, which will be discussed in the next section.

3.6 Constraints

Given the multitude of different astrophysical equations of state available, it's necessary to limit those implemented in our supernova simulations to those which are most relevant in light of the most recent observational and experimental data. Advancements in astrophysical observation in the past few decades provide neutron star observations that allow the mass and radius of the star to be measured with unprecedented precision, constraining the mass–radius phase space allowed for neutron stars (Demorest et al. 2010; Antoniadis et al. 2013). Precision in the experimental measurement of nuclear masses, binding energies, and resonances has also improved, allowing for improved characterization and constraint of the nuclear parameters in any given equation of state. Each constraint can limit the viability of a given equation of state in a different region of the phase space in temperature, density, or composition. This is a limitation of the constraints, in that the constraints are often only applicable in a limited range of the phase space. Most laboratory experiments test the saturation properties of symmetric nuclear matter, unlike the extremely neutron-rich conditions found in the PNS. Laboratory experiments can't yet reproduce the full range and extent of the extreme conditions in density, temperature, and composition found in CCSNe. Astrophysical observations are of colder, denser neutron stars than the conditions found in much of a supernova simulation. For regions not covered by these constraints, constructing an equation of state requires the use of theoretical calculations and extrapolations to fill the phase space required in order to simulate the CCSN event (see, e.g. Oertel et al. (2017)).

3.6.1 Theoretical Neutron Matter Calculations

In contrast to the theoretical description of the nuclear interactions for symmetric matter, pure neutron matter has a simpler nuclear interaction Hamiltonian due to its simplified isospin structure. Because of this, theoretical descriptions of neutron matter can advance the ability to constrain equation of state models as they approach dense neutron matter. At densities up to nuclear matter saturation, many different theoretical methods can be used to calculate the nuclear interaction Hamiltonian, specifically with regards to two- and three-body neutron interactions. Some (but not all) of these nuclear treatment methods are as follows: phenomenological Skyrme force methods (Stone and Reinhard 2007), variational calculations using nuclear potential models (Friedman et al. 1981), Brueckner-Hartree-Fock methods used to investigate asymmetric non-relativistic nuclear matter (Baldo et al. 1998; Vidaña et al. 2011), Quantum Monte Carlo (QMC) methods, such as Auxiliary Field Diffusion Monte Carlo (AFDMC) (Lattimer and Lim 2013; Gandolfi et al. 2010; Gandolfi 2013), and more recently chiral effective field theory (χEFT) (Hebeler et al. 2010; Tews et al. 2013). The inclusion of three-nucleon forces adds significant repulsion up to and at nuclear saturation density. Additional bounds on the symmetry energy J and its slope parameter L have been put forward by comparison of the energy of unitary gas models to the energy of pure neutron matter, which in turn constrains the incompressibility (Tews et al. 2017).

3.6.2 Nuclear Physics Laboratory Measurements

One of the primary constraints from laboratory experiments is the fitting from experimentally determined nuclear masses and binding energies, as this method is largely model independent and the results from such experiments are readily available and of high accuracy. Extrapolation techniques can be used to investigate nuclear matter properties for exotic heavier nuclei. Comparison with Hartree-Fock and Thomas-Fermi models, as well as FRDM, make it possible to glean constraints on the symmetry energy J and its slope parameter L resulting from those models when fitted to the nuclear mass data. This has shown a linear correlation

between the symmetry energy at nuclear saturation and L that has been experimentally verified using multiple approaches (Lattimer and Steiner 2014; Nazarewicz et al. 2014). Additionally, constraints on the nucleon-nucleon interactions in a given equation of state can be gleaned from experimentally determined giant and pygmy nuclear dipole resonances (Daoutidis and Goriely 2011; Fattoyev et al. 2010).

Extensive investigation of the neutron skin thickness, r_{np} , of the neutron-rich nucleus ^{208}Pb has shown that r_{np} is a function of both the symmetry energy (specifically, the volume component of the symmetry energy from the FRDM formulation), and the slope parameter L (Lattimer and Lim 2013). The resulting relationship is that the faster the symmetry energy increases with density, the larger the size of the neutron skin in heavy nuclei (Roca-Maza et al. 2011). Additional neutron skin thickness measurement programs are underway via the CREX and PREX experiments at Jefferson Lab (Horowitz et al. 2014).

Heavy-ion collisions are an additional source of constraint on the EOS from laboratory experiments, as they are a probe of densities several times greater than nuclear saturation density (albeit over the short time of the beam collision), and of temperatures of 10-50 MeV, conditions similar to what is found in CCSNe. Additionally, constraints on the stiffness of cold nuclear matter can be estimated from transverse and elliptic flow observables in heavy ion collisions. However, the matter in heavy ion collisions is not usually as isospin asymmetric as can be found in the high density regions of the CCSNe. Likewise, the timescales in heavy-ion collisions are much smaller than in the dense matter regime of the PNS with respect to weak equilibrium; these differences must be accounted for when using heavy-ion collision data for constraints (Tsang et al. 2009).

3.6.3 Astrophysical Observations

The primary constraints from astrophysical observation have come from high-precision measurements of binary systems consisting of a neutron star and a white dwarf, or pairs of neutron stars. Only in the last two decades have these observations become of sufficient

precision to conclusively characterize the masses and radii of the constituent compact stars, as seen in Figure 3.1. In one instance, the mass of a pulsar, a rapidly spinning neutron star, was determined by examination of the general relativistic Shapiro delay of the pulsar’s signal as it passed by its white dwarf companion (Demorest et al. 2010), which led to a neutron star mass measurement of $1.928 \pm 0.017 M_{\odot}$. Additionally, a second instance of a neutron star/white dwarf binary in which analysis of its orbital decay combined with high confidence in the mass of the white dwarf companion measured a neutron star mass of $2.01 \pm 0.01 M_{\odot}$ (Antoniadis et al. 2013). Most useful to constraining the equation of state from astrophysical data would be a accurate determination of both the mass and radius for the same neutron star. This was first achieved via observations of X-ray bursts from photospheric radial expansions, and also via observations of thermal emissions from the neutron star surface in quiescent low-mass X-ray binaries, allowing for the inference of the mass and radii simultaneously via the angular emission area (Steiner et al. 2010). However the mass constraint provided by this method is much coarser than the Shapiro delay observations.

The general relativistic Tolman-Oppenheimer-Volkov (TOV) equation determines the stellar structure of the neutron star, and thus with the mass and radius observations, the TOV equation can be integrated in order to quantify the relationship between the pressure and density in the nuclear matter equation of state (Tolman 1939; Oppenheimer and Volkoff 1939),

$$\frac{dP}{dr} = -\frac{G \left[m(r) + 4\pi r^3 \frac{P}{c^2} \right] \left[\epsilon + \frac{P}{c^2} \right]}{r \left[r - \frac{2Gm(r)}{c^2} \right]} \quad (3.8)$$

$$\frac{dm(r)}{dr} = 4\pi \epsilon r^2 \quad (3.9)$$

where the pressure is $P = \rho^2 \left(\frac{\partial E}{\partial \rho} \right)$, the energy density is $\rho(E + m_n)$, G is Newton’s gravitational constant, $m(r)$ is the enclosed gravitational mass at radius r , and m_n is the mass of the neutron. The integration of these equations gives the profiles of enclosed mass, pressure, energy density, and density as functions of radius. The key constraint from

measured masses is that a particular EOS must provide for the existence of a neutron star of at least this mass. Observations of neutron stars of $\sim 2 M_{\odot}$ imply that any EOS in which the neutron star collapses to a black hole at less than the observed mass can therefore not be an accurate description of these astrophysical and nuclear matter conditions.

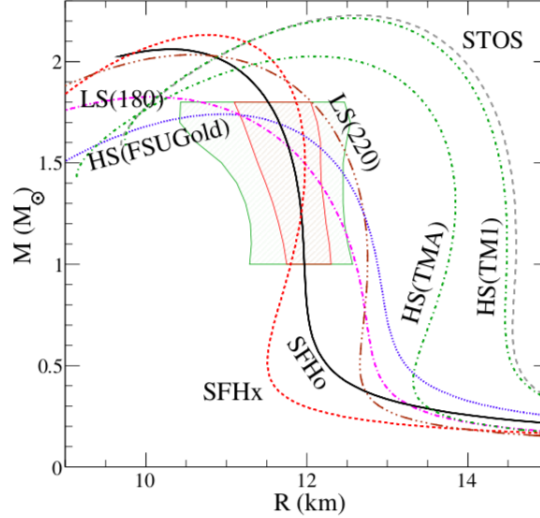


Figure 3.1: Constraint of Neutron Star Mass vs. Radius, by EOS. (Steiner et al. 2013)

The combined result of these experimental, observational, and theoretical approaches allows for the constraint of the incompressibility K to 240 ± 20 MeV (Piekarewicz 2010; Steiner et al. 2013). Similarly, the symmetry energy has been constrained to $28 \lesssim E_{\text{sym}} \lesssim 34$ MeV (Tsang et al. 2009). The density dependent slope parameter L is less well constrained due to model-dependent interpretation of experiments, putting it in the range of $20 \lesssim L \lesssim 120$ MeV (Carbone et al. 2010).

3.7 Overview of Relevant Equations of State

For our purposes, we will investigate seven particular astrophysical/supernova equations of state. The default equation of state for the past two decades of CCSN simulation has been that of Lattimer & Swesty (LS) (Lattimer and Swesty 1991). Formerly in widespread use

was the equation of state of H. Shen, Toki, Oyamatsu & Sumiyoshi (STOS) (Shen et al. 1998). In more recent years, additional equations of state have become available, as discussed in the coming sections.

3.7.1 Lattimer & Swesty

The LS equation of state uses a compressible, non-relativistic liquid drop model to describe nuclear matter and nuclei together with dripped nucleons. A Skyrme-type effective interaction force model is used to calculate nuclear many-body interactions, and from this the density dependence of the equation of state is determined (Lattimer and Swesty 1991). The Maxwell construction (a method for maintaining thermodynamic consistency) is used to connect the nuclei and the nuclear matter regimes. Electrons are treated as non-interacting ultra-relativistic particles in pair equilibrium. Alpha particles are assumed to represent the distribution of light nuclei, while the SNA is implemented to reflect an ensemble of heavy nuclei. The Wigner-Seitz approximation is used to treat heavy ion interaction. Nuclear shell effects are not included. Bulk free energy treatment is the same for nucleons and nuclei. The Thomas-Fermi model is used for surface corrections. For LS, a nuclear saturation density of $\rho_s = 2.59 \times 10^{14} \text{ g cm}^{-3}$ and a symmetry energy $E_{\text{sym}} = 29.3 \text{ MeV}$ are assumed. Historically, the LS equation of state has been seen as the limiting “soft” equation of state in CCSNe simulations. In this study, we use the $K = 220 \text{ MeV}$ incompressibility version of the LS EoS for $\rho > 10^{11} \text{ g cm}^{-3}$, as it is the version of LS that meets the most stringent constraints.

3.7.2 Baron, Cooperstein, & Kahana

In the simulation utilizing LS, for $\rho < 10^{11} \text{ g cm}^{-3}$ Chimera utilizes an enhanced version of the Baron, Cooperstein, and Kahana (Baron et al. 1985) EoS. This is due to a known underestimation in the LS EoS in the alpha mass fraction at these densities. Additionally, the Cooperstein electron–positron EoS is used for all simulations.

3.7.3 STOS

In contrast to the LS equation of state, the Shen, Toki, Omayatsu, Sumiyoshi (STOS) equation of state does not use a liquid drop model, but instead introduced the use of a Relativistic Mean Field (RMF) model to handle many-body interactions, not by summing up the mutual two-body interactions of the constituent particles, but instead to describe the interaction of one particle with the remaining ones by an average potential created by the others (Shen et al. 1998). The specific parameters of this RMF model are those based on the TM1 RMF interaction (Sugahara and Toki 1994). In the TM1 interaction, the parameter set was fitted to the binding energies and charge radii of heavy nuclei. Like the LS equation of state, electrons are assumed to be non-interacting, so LS electron and photon contributions are used for continuity. The Thomas-Fermi and variational approximations are used to describe non-uniform nuclear matter and nuclei. The SNA is also assumed, and nuclear shell effects are not included. Historically, the STOS equation of state has been the limiting “stiff” case in CCSNe simulations (Suwa et al. 2013).

3.7.4 DD2

The DD2 EOS uses a quantum statistical approach with a density-dependent (hence the “DD” of DD2) RMF for nuclear interactions. The density dependence arises from replacing the regular meson scalar and vector self-energies with density dependent forms. The “2” in DD2 represents the improvement over the original DD nuclear parameter set via the incorporation of newer experimental nucleon mass data (Typel 2005). For the DD2 EOS, binding energies for nuclei without experimentally measured values are calculated with the FRDM. The ensemble of nuclei is represented by the HS treatment.

3.7.5 Steiner, Hempel, & Fischer

For SFHo, the nuclear matter equation of state is constrained by mass and radius measurements from quiescent low-mass X-ray binaries and neutron star observations (Steiner et al. 2010). Such observations suggest smaller neutron star radii and lower pressures just above nuclear saturation density, which in turn implies more gravitational binding energy relative to previous assumptions. The strong force interaction for nuclear matter is parametrized with a new Walecka-model-based RMF model that includes additional fitting via binding energy measurements and charge radii measurements for ^{208}Pb and ^{90}Zr . The density dependence of the symmetry energy is relatively flexible due to the inclusion of additional non-linear couplings and self-couplings in the interaction Lagrangian (Steiner et al. 2005, 2013). SFHx is a variant that specifically minimizes the slope of the symmetry energy L for neutron star compactness (hence the “x” designation, for extreme, versus the “o” for ordinary). The binding energy table is the same as is used for DD2 and the ensemble of nuclei is represented by the HS treatment.

3.7.6 FSUGold & IUFSU

The FSUGold EOS implements the nuclear interaction using an RMF with additional omega- and rho-meson couplings in the Lagrangian, which affects the density dependence of the symmetry energy (Todd-Rutel and Piekarewicz 2005; Piekarewicz 2007). The IUFSU EOS was developed using FSUGold as a starting point by including more recent experimental and observational data (Fattoyev et al. 2010). IUFSU includes additional fitting from experimental investigation of neutron skin thickness of ^{208}Pb , as well as further fitting from the observationally deduced mass of the $2.01 M_{\odot}$ neutron star mentioned earlier (Antoniadis et al. 2013). The binding energies for FSUGold and IUFSU come from the table of X. Roca-Maza, which contains 1512 species (Roca-Maza and Piekarewicz 2008). The ensemble of nuclei is represented by the HS treatment.

3.7.7 NL3

The NL3 EOS utilizes an RMF fitting procedure using the binding energies, neutron radii, and charge radii of 10 specific nuclei, ^{16}O , ^{40}Ca , ^{58}Ni , ^{90}Zr , ^{116}Sn , ^{124}Sn , ^{132}Sn , ^{208}Pb , and ^{214}Pb (Lalazissis et al. 1999). The Lagrangian includes cubic and quartic scalar self-interactions, and the binding energy table contains 1315 species (Lalazissis et al. 1999). The ensemble of nuclei is represented by the HS treatment.

3.8 Choices of EOS for Comparison

For the purposes of this study, EOS's which violate these the most modern constraints were excluded, including the historically relevant STOS EOS. The unitary gas boundary on symmetry energy parameters excluded both EOS's as their RMF choices, TM1 and TMA, were both excluded, as well as more recently LS220 (Tews et al. 2017). NL3 was also excluded by this constraint. We nonetheless chose to include it in the simulation set due to the fact that it produces the largest neutron star mass and this extremal case allows for the better understanding of the ranges of effects of the EOS. Additionally, astrophysical constraints rule out the STOS EOS, but not NL3 (Steiner et al. 2013). Other than LS220 and NL3, the EOSs listed in Table 3.1 meet the constraints. In this comparison, the RMF group refers to all EOS other than LSBCK, as all EOS discussed here other than LSBCK utilize an RMF approximation.

Table 3.1: Summary of EOS Nuclear Parameters

	SFHo	SFHx	DD2	IUFSU	FSUGold	NL3	LSBCK
$\rho_s(\times 10^{14} \text{ g cm}^{-3})$	2.63	2.66	2.48	2.57	2.46	2.46	2.57
E_0 (MeV)	16.19	16.16	16.02	16.39	16.27	16.24	16.0
K (MeV)	245.4	238.8	242.7	231.3	229.5	271.5	220
K' (MeV)	-467.8	-457.2	168.7	-290.3	-523.9	202.6	243.2
J (MeV)	31.57	28.67	31.67	31.29	32.56	37.39	29.3
L (MeV)	47.10	23.18	55.03	47.20	60.43	118.49	74
K_{sym} (MeV)	-205.4	-40.0	-93.2	28.5	-51.4	100.8	64.9
$M_{max} M_\odot$	2.06	2.13	2.42	1.95	1.74	2.79	2.06
$R_{M_{max}}$ (km)	10.32	10.77	11.90	11.31	10.95	13.40	10.67
$R_{1.4}$ (km)	11.9	12.0	13.2	12.7	12.6	14.8	12.71

Chapter 4

Numerical Methods & Physics Inputs

In this section I will firstly describe the WeakLib infrastructure, developed to allow for the CHIMERA supernova simulation code to utilize multiple modern equations of state in a plug-and-play fashion. I will discuss the development, capability, and testing of this EOS framework. I will then discuss the numerical implementation of the EOS physics in CHIMERA as well as the other aspects of the CHIMERA code.

4.1 WeakLib

Sophisticated implementation of both the EOS as well as the neutrino microphysics and transport are necessary to accurately model the CCSN mechanism given the roles of neutrinos in reviving the shock, PNS neutronization, and cooling. The evolution of the neutrino radiation field is characterized by the general relativistic Boltzmann kinetic equation. The long-term goal of WeakLib is to enhance the computational efficiency of the neutrino microphysics implementation in CCSN simulations via pre-processed state-of-the-art neutrino interaction opacity tables for use in the collision term (currently in development by my fellow graduate student Ran Chu), which would be usable in a straightforward manner by CHIMERA, but specific discussion of the neutrino interactions is beyond the scope of

this dissertation. However, given the sensitivity of neutrino microphysics interactions to the EOS, it behooves us to first create the capacity for our simulations to use the latest EOS's available, of which there are many, dependent on different nuclear force models, as discussed previously. WeakLib is the framework via which these varied EOS's have been put into tabular form compatible with use by CHIMERA (as well as other codes currently in development). WeakLib uses pre-tabulated EOS data from the Compstar Online Supernova Equations of State repository (CompOSE).¹

4.1.1 Implementation

The organization of the WeakLib framework is shown in Figure 4.1. It is firstly a repository for the subroutines to handle the necessary input/output (I/O) for reading and writing the EOS tables, as well as the subroutines and modules for handling table operations for direct use of the tables in applications. All I/O is done using the hierarchical data format (HDF) HDF5 I/O library. I will discuss the individual sections of the code structures, and what they contain.

Distribution The distribution section of the framework is what is meant to be distributed to users of WeakLib for use directly or for the purpose of implementing WeakLib into simulation codes. This section contains the subroutines to build, read, write, and utilize the tables.

EOSSource

EOSSource contains EquationOfStateModule subroutines for loading, allocating, deallocating, and handling EOSs.

¹compose.obspm.fr

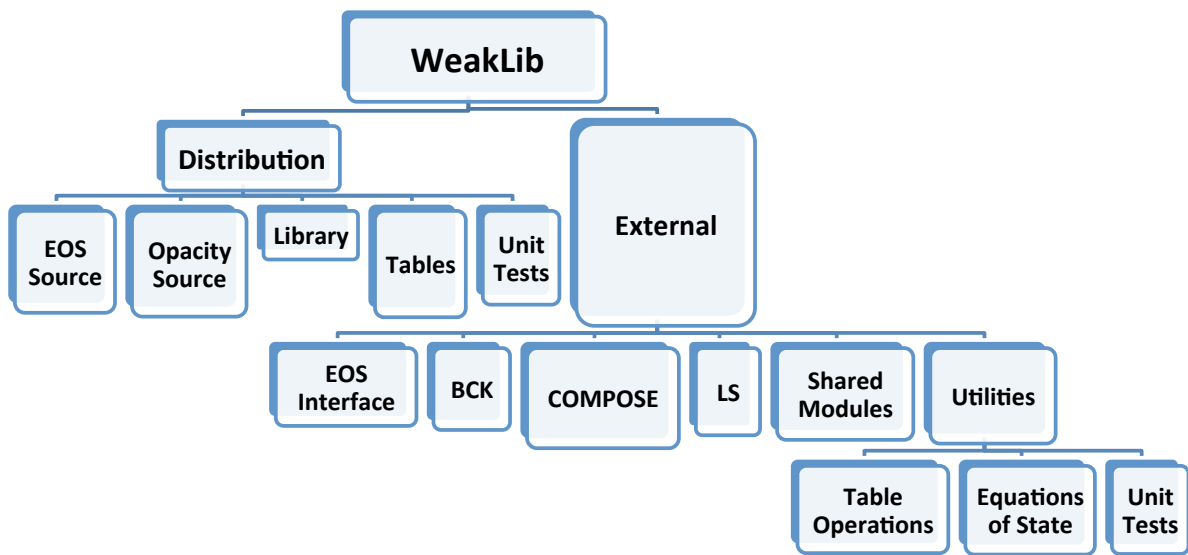


Figure 4.1: WeakLib Code Structure

OpacitySource

OpacitySource is the same as EOSSource, but for opacities.

UnitTests

UnitTests includes tests for Library modules.

Library

Library contains modules for interpolation & derivation, grid-making, I/O, parallelization, ThermoState, DependentVariables, and associated subroutines.

4.1.2 External

The external section of the framework is specifically for WeakLib code development and testing. This is where the LS table, known to have bad values in certain extremes of the phase space it covers, was repaired via interpolation between valid data points.

LS

This section contains the LS EOS data.

BCK

This section contains the BCK EOS data.

EOS Interface

This section contains the subroutines for accessing the data from the LS and BCK EOSs and combining them as needed.

Shared Modules

This section contains modules for storing physical constants and numerical routines used by both individual unit tests and the LS and BCK EOSs.

COMPOSE

The COMPOSE section contains the subroutines for building raw data tables from CompStar full data tables for each EOS. The CompStar tables for each EOS contain more data than is needed for the principle applications of WeakLib, but raw tables can be constructed with whatever subset of the CompStar provided data needed.

Utilities

Utilities contains modules for building, restructuring, repairing, and manipulating the table structures.

Equations of State

This is the repository that contains the drivers for creating tables for a give equation of state created in the COMPOSE section and moving the data into a more universal format for use in application. The CompStar tables do not contain the binding energies, so this repository also contains the binding energy tables supplied by Matthias Hempel² and the subroutines needed to access them.

Table Operations

Contains modules for characterizing holes in the table (where a Newton-Raphson approximation iteration fails or a NaN occurs), filling them via interpolation, and checking for

²<http://phys-merger.physik.unibas.ch/~hempel/eos.html>

monotonicity.

Unit Tests

This section contains all unit tests that are specific to a particular subroutine being tested or EOS being constructed. The raw data extracted from the CompStar EOS data table is read and put into WeakLib standard format here.

4.1.3 The Tables

The table structure consists of HDF5 files that are read into a FORTRAN derived datatype containing additional derived datatypes, one storing the thermodynamic state of intrinsic variables of density, temperature, and electron fraction (called ThermoState), and another containing the dependent variables (called Dependent Variables). The tables are put into a format such that for a given 3-tuple of the intrinsic variables of density, temperature, and electron fraction, all of the corresponding dependent variables can be obtained:

1. Pressure (dynes per cm^2)
2. Entropy per baryon (k_b per baryon)
3. Internal energy density (erg per gram)
4. Electron chemical potential (MeV)
5. Proton chemical potential (MeV)
6. Neutron chemical potential (MeV)
7. Proton Mass Fraction
8. Neutron Mass Fraction

9. Alpha Mass Fraction
10. Heavy Mass Fraction
11. Heavy Charge Number
12. Heavy Binding Energy (MeV)
13. Thermal Energy (MeV)
14. Gamma1

where Gamma1 is the first adiabatic coefficient.

4.1.4 Table Grid

Given that the EOS significantly influences the evolution of many sectors of the supernova explosion dynamics, the dependent variables of the EOS which influence these dynamics must be recalculated or re-obtained as the state variables of the fluid change, such that the fidelity and thermodynamic consistency of the supernovae simulation is preserved. In modern multi-dimensional supernova codes, this occurs several times in a given cycle; the nuclear burning network, the hydrodynamic step, and the transport step all query the EOS for the values of the dependent variables (pressure, internal energy, entropy, chemical potentials, mass fractions, etc.) and in some cases for their derivatives with respect to the state variables. In order that the EOS data returned by these queries is of the highest fidelity possible, we interpolate the data on a grid in density, temperature, and electron fraction of a chosen resolution and range.

For WeakLib, I have constructed tabular EoS tables, starting with a 3-dimensional array of logarithmically increasing density and temperature, and linearly increasing electron fraction. At the conditions which occur inside this region of the star, the thermodynamic state variables can be determined exactly from the density, temperature, and electron fraction. Using the vertices of this array, the Lattimer & Swesty (220) EoS is called, and the dependent

state variables are populated. For each of the CompStar tables, I specified the grid required in their driver subroutines, and then populated the tables into the HDF5 format used by CHIMERA.

4.1.5 Interpolation Methods

In order to provide accurate state variable values for (ρ, T, Y_e) 3-tuples that fall between the table vertices, we implement a tri-linear interpolation method. Firstly, we locate the ρ , T , and Y_e values above and below the triplet values,

$$\Delta_1 = \frac{\log_{10}\left(\frac{\rho}{\rho_1}\right)}{\log_{10}\left(\frac{\rho_2}{\rho_1}\right)}; \Delta_2 = \frac{\log_{10}\left(\frac{T}{T_1}\right)}{\log_{10}\left(\frac{T_2}{T_1}\right)}; \Delta_3 = \frac{(Y_e - Y_{e1})}{(Y_{e2} - Y_{e1})} \quad (4.1)$$

then use first-order tri-linear interpolation as follows, for a given dependent variable v :

$$\begin{aligned} (\bar{v}) = & (1 - \Delta_3)((1 - \Delta_1)(1 - \Delta_2)\bar{v}(\rho_1, T_1, Y_{e1}) \\ & + \Delta_1(1 - \Delta_2)\bar{v}(\rho_2, T_1, Y_{e1}) \\ & + (1 - \Delta_1)\Delta_2\bar{v}(\rho_1, T_2, Y_{e1}) \\ & + \Delta_1\Delta_2\bar{v}(\rho_2, T_2, Y_{e1})) \\ & + \Delta_3((1 - \Delta_1)(1 - \Delta_2)\bar{v}(\rho_1, T_1, Y_{e2}) \\ & + \Delta_1(1 - \Delta_2)\bar{v}(\rho_2, T_1, Y_{e2}) \\ & + (1 - \Delta_1)\Delta_2\bar{v}(\rho_1, T_2, Y_{e2}) \\ & + \Delta_1\Delta_2\bar{v}(\rho_2, T_2, Y_{e2})) \end{aligned} \quad (4.2)$$

$$\begin{aligned}
\left(\frac{dv}{d\rho}\right)_{T,Y_e} &= v\alpha_1[(1-\Delta_3)[-(1-\Delta_2)\bar{v}(\rho_1, T_1, Y_{e1}) \\
&\quad + (1-\Delta_2)\bar{v}(\rho_2, T_1, Y_{e1}) \\
&\quad - \Delta_2\bar{v}(\rho_1, T_2, Y_{e1}) \\
&\quad + \Delta_2\bar{v}(\rho_2, T_2, Y_{e1})] \\
&\quad + \Delta_3[-(1-\Delta_2)\bar{v}(\rho_1, T_1, Y_{e2}) \\
&\quad + (1-\Delta_2)\bar{v}(\rho_2, T_1, Y_{e2}) \\
&\quad - \Delta_2\bar{v}(\rho_1, T_2, Y_{e2}) \\
&\quad + \Delta_2\bar{v}(\rho_2, T_2, Y_{e2})]]
\end{aligned} \tag{4.3}$$

$$\begin{aligned}
\left(\frac{dv}{dT}\right)_{\rho,Y_e} &= v\alpha_2[(1-\Delta_3)[(1-\Delta_1)\bar{v}(\rho_1, T_1, Y_{e1}) \\
&\quad + \Delta_1\bar{v}(\rho_2, T_1, Y_{e1}) \\
&\quad - (1-\Delta_1)\bar{v}(\rho_1, T_2, Y_{e1}) \\
&\quad - \Delta_1\bar{v}(\rho_2, T_2, Y_{e1})] \\
&\quad + \Delta_3[(1-\Delta_1)\bar{v}(\rho_1, T_1, Y_{e2}) \\
&\quad + \Delta_1\bar{v}(\rho_2, T_1, Y_{e2}) \\
&\quad - (1-\Delta_1)\bar{v}(\rho_1, T_2, Y_{e2}) \\
&\quad - \Delta_1\bar{v}(\rho_2, T_2, Y_{e2})]]]
\end{aligned} \tag{4.4}$$

$$\begin{aligned}
\left(\frac{dv}{dY_e}\right)_{\rho,T} &= v\alpha_3[-(1-\Delta_1)[(1-\Delta_2)\bar{v}(\rho_1, T_1, Y_{e1}) \\
&- \Delta_1(1-\Delta_2)\bar{v}(\rho_2, T_1, Y_{e1}) \\
&- (1-\Delta_1)\Delta_2\bar{v}(\rho_1, T_2, Y_{e1}) \\
&- \Delta_1\Delta_2\bar{v}(\rho_2, T_2, Y_{e1})] \\
&+ (1-\Delta_1)[(1-\Delta_2)\bar{v}(\rho_1, T_1, Y_{e2}) \\
&+ \Delta_1(1-\Delta_2)\bar{v}(\rho_2, T_1, Y_{e2}) \\
&+ (1-\Delta_1)\Delta_2\bar{v}(\rho_1, T_2, Y_{e2}) \\
&+ \Delta_1\Delta_2\bar{v}(\rho_2, T_2, Y_{e2})]
\end{aligned} \tag{4.5}$$

where

$$\alpha_1 = \frac{1}{\log_{10}\left(\frac{\rho_2}{\rho_1}\right)}; \alpha_2 = \frac{1}{\log_{10}\left(\frac{T_1}{T_2}\right)}; \alpha_3 = \frac{1}{(Y_{e2} - Y_{e1})} \tag{4.6}$$

4.1.6 Test Problems

In order to best understand the sensitivity of the the EOS dependent variables, we here show how the doubling of the resolution in ρ, T , and Y_e independently affects the accuracy of our interpolation vs. the direct call to the CHIMERA default LS EOS, comparing how the doubling of resolution in ρ and T from 20 points per decade to 40 points per decade effects the relative error,

$$\delta = \left| \frac{v - v_{\text{interpolated}}}{v} \right| \tag{4.7}$$

where v is a given dependent variable. The following Figures 4.2 to 4.8 involve the data from the direct query of the LS EOS for a given (ρ, T, Y_e) 3-tuple versus the query of a WeakLib constructed LS table for the variable value for the same 3-tuple using interpolation between the table vertices.

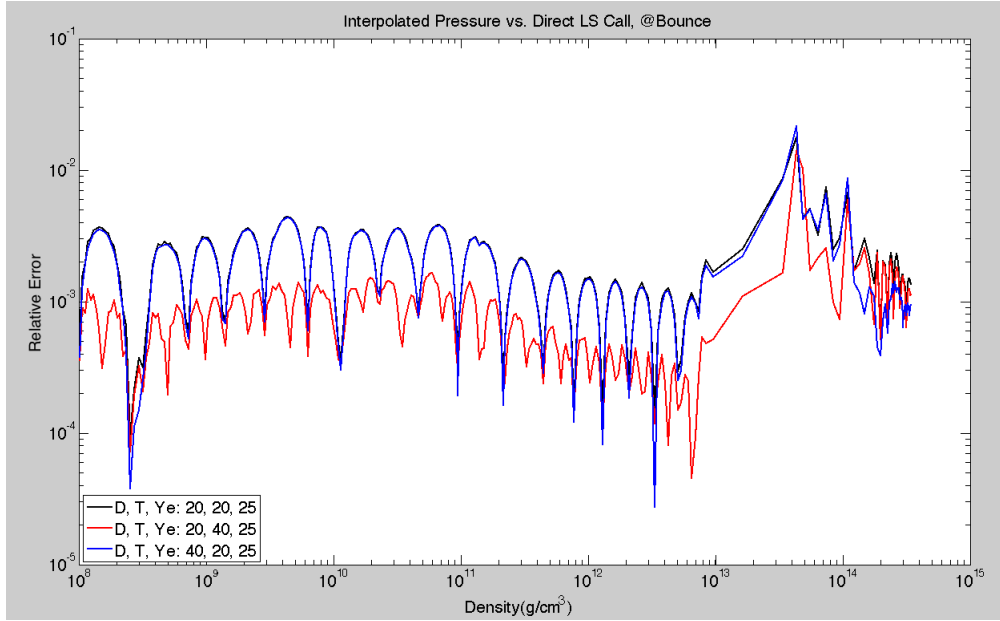


Figure 4.2: Relative pressure error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.

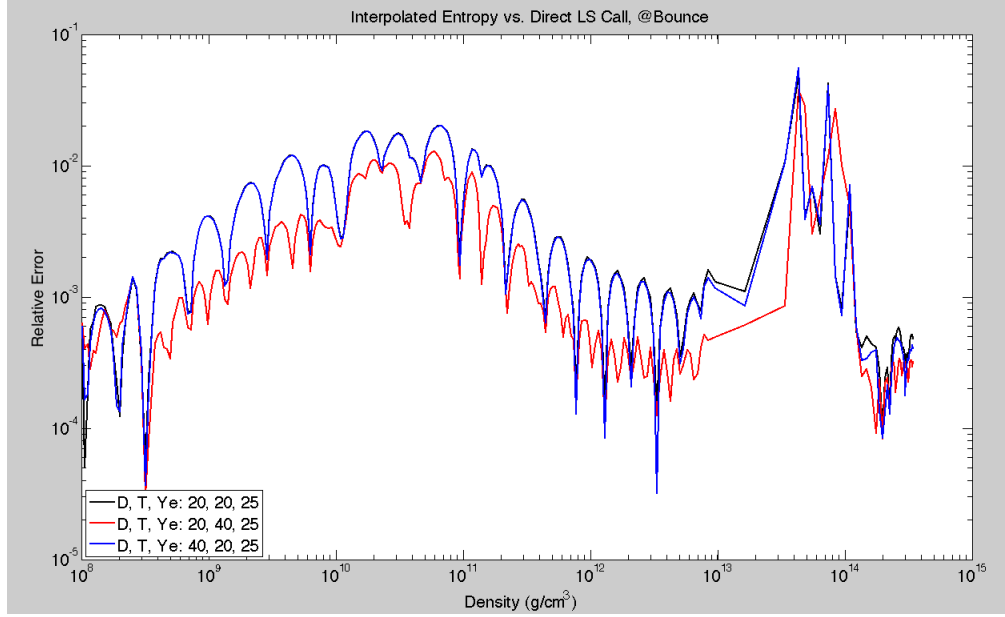


Figure 4.3: Relative entropy error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.

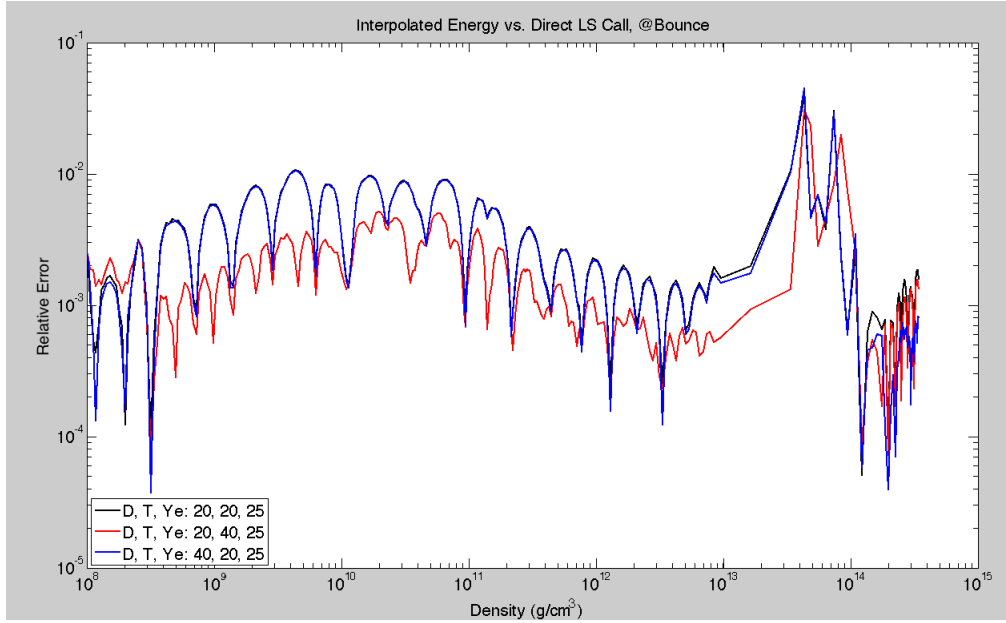


Figure 4.4: Relative internal energy error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.

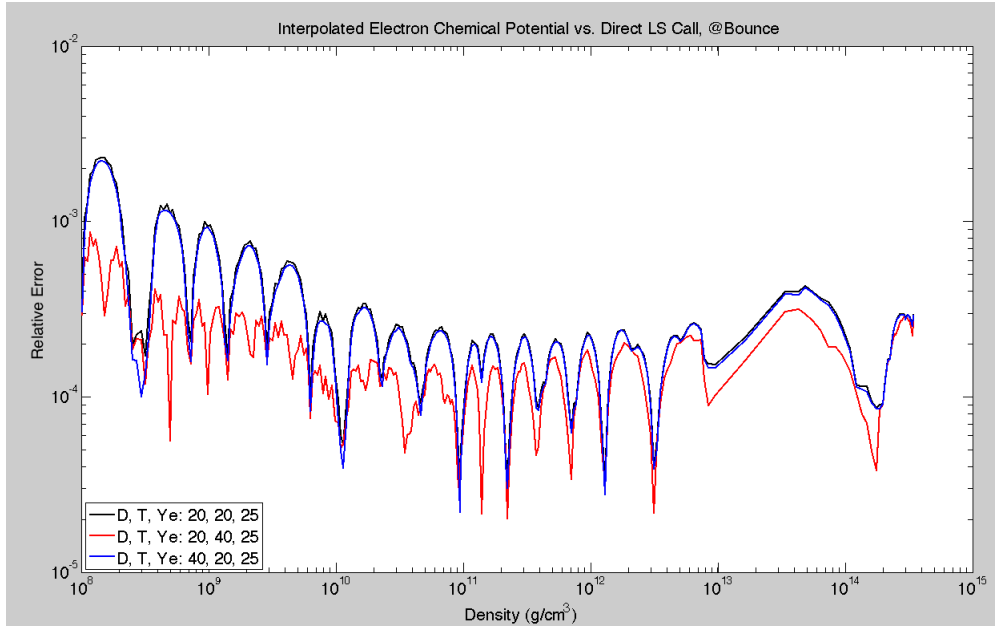


Figure 4.5: Relative electron chemical potential error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.

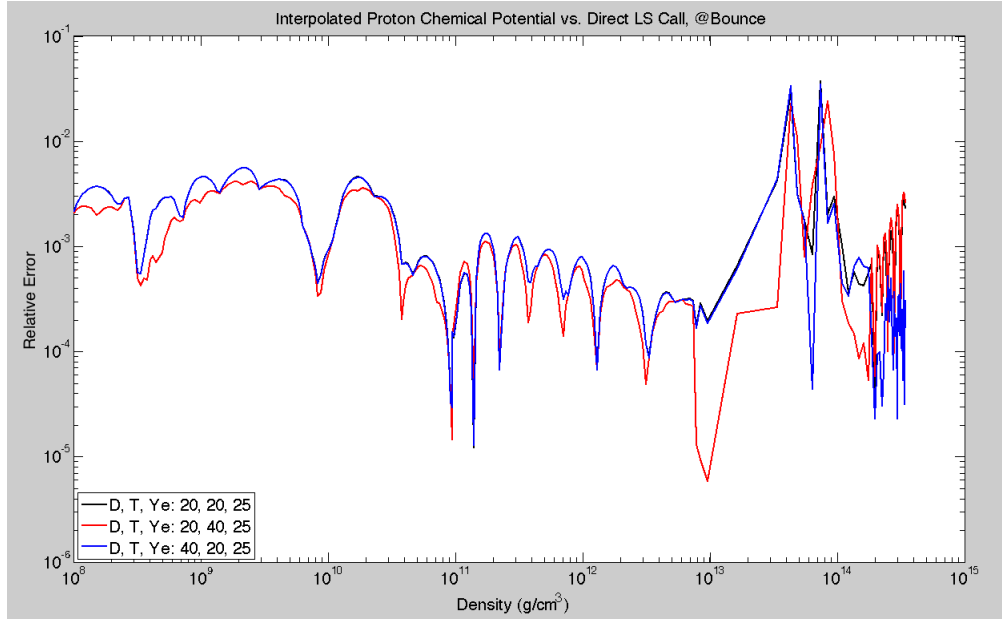


Figure 4.6: Relative electron chemical potential error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.

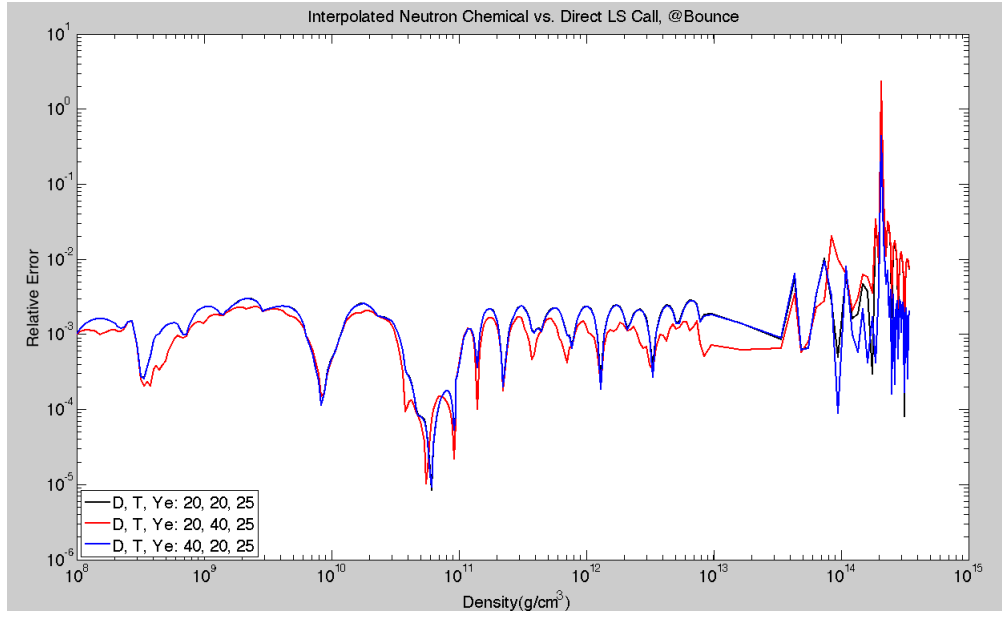


Figure 4.7: Relative electron chemical potential error between direct output from the LS EOS versus interpolated quantities for specified table resolutions.

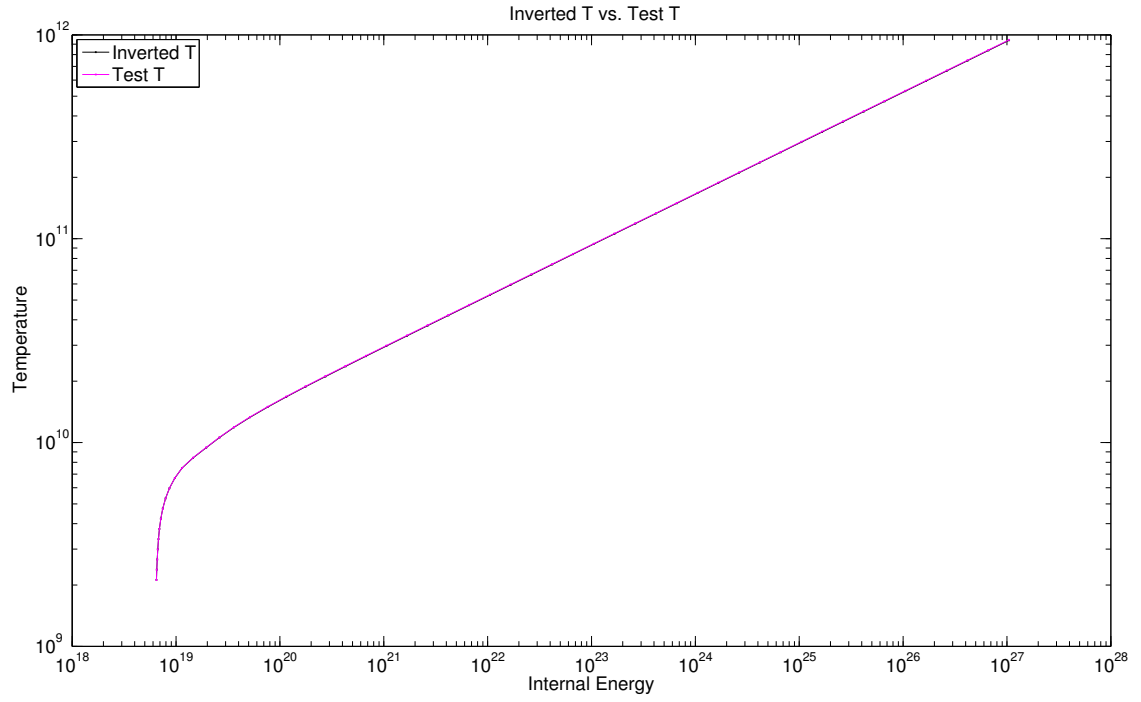


Figure 4.8: Inversion Errors. This graph shows the correlation between temperature and internal energy of the LS EOS when queried from the table directly versus when the temperature is inversely calculated from the internal energy.

4.1.7 Thermodynamic Consistency and Self Consistency

The EOSs from the CompStar database are constructed with inbuilt thermodynamic consistency. Though the values of the mass fractions pre-interpolation equals one, linear interpolation of non-linear variables may lead to the errors post-interpolation in the sum of the mass fractions. In order to ensure thermodynamic consistency, abundance distribution is normalized to insure that mass fractions of the constituent particles always sums to one when the table is put into WeakLib format:

$$X_p + X_n + X_\alpha + \sum_h X_h = 1 \quad (4.8)$$

This was in particular necessary several of the tables presented here (such as SFHo and SFHx) contain light nuclear clusters (such as ^3He , ^3H , and ^2H), which are divided into their constituent protons and neutrons and combined with those quantities. Additionally, the LS table required repair in certain regions where its construction via Newton-Raphson methods fails, so subroutines for hole identification and repair were developed and implemented, but that is only necessary for the LS EOS.

4.1.8 Extrapolation Methods

In the course of our simulations, the electron fraction of the conditions in the supernova went out of the range of each table, firstly for the softer equations of state but eventually for all of them. In order to explore the physics of this high electron fraction regime, we used quadratic extrapolation where possible (high temperature), and where it lead to inconsistent values, we recalculated the mass fractions using the Saha equation, Equation [3.1](#).

4.2 CHIMERA

To model the effects of the EOS on the supernova mechanism, I have used the CHIMERA³ code, a parallel, multi-physics, multi-dimensional supernovae simulation code, named after the eponymous mythical creature. Like the creature, the CHIMERA code has three principle “heads”: a spectral neutrino transport code, a shock-capturing hydrodynamics code, and a nuclear kinetics code (Bruenn et al. 2018, in preparation). In addition, it contains components for the incorporation of self-gravity, for the nuclear equation of state, and a series of overlayers for simulation management and incorporating each component into the whole, as well as the handling of I/O. Additionally, each CHIMERA simulation creates additional output data for the post-processing so as to give accurate neutrino & gravitational wave signatures, as well as output of nucleosynthesis during the supernova event. I will discuss each of these parts of the overall code, with a particular eye to how they interplay with the nuclear EOS.

4.2.1 Hydrodynamics

The hydrodynamics are evolved via a dimensionally-split Lagrangian-remap scheme using an modified form of the piecewise parabolic method (PPMLR; Colella and Woodward (1984) as originally implemented in the multidimensional ideal compressible hydrodynamics code VH1 (Hawley et al. (2012)). In order to reflect the extremely high densities that occur during CCSNe and the formation of the PNS, our simulations include the effects of General Relativity (GR) via the use of multipole Newtonian self-gravity with spherical GR corrections. As the equation of state governs the relationship between thermodynamic quantities such as pressure, entropy, internal energy, and temperature, it plays into the hydrodynamics by providing closure to the hydrodynamic equations.

³ChimeraSN.org

4.2.2 Neutrino Transport/Radiation Hydrodynamics

Neutrino transport is modeled using a ray-by-ray plus (RbR+) approximation (Buras et al. 2003), the plus reflecting the inclusion of lateral neutrino pressure gradients. The transport solver is an improved and updated version of the multi-group, flux-limited, diffusion transport solver of Bruenn (1985), enhanced for GR (Bruenn et al. 2001). This method computes and independent solve along each radial ray, and includes a geometric flux limiter so as to prevent non-physical behavior, specifically limiting spatial derivatives to realistic values. This treatment includes four neutrino species, electron neutrinos and their antiparticles, and μ and τ neutrinos and their commensurate antiparticles, using 20 logarithmically spaced energy groups (4 – 250 MeV). Neutrinos are advected laterally (in the θ -direction) with the fluid and contribute to the lateral pressure gradient where $\rho > 1 \times 10^{12} \text{ g cm}^{-3}$. During evolution the radial zones are gradually and automatically repositioned during the remap step to track changes in the radial structure. To minimize restrictions on the time step from the Courant limit, the lateral (θ) hydrodynamics for the inner $0.5 M_{\odot}$ are frozen during collapse, and, after prompt convection fades, we expand the laterally frozen region to the inner 6 km. In the “frozen” region we set $v_{\theta} = 0$ and skip the lateral hydrodynamic sweep. The full radial hydrodynamics and neutrino transport are always computed to the center of the simulation for all angular rays.

4.2.3 Weak Physics Inputs

The effectiveness of neutrino heating for a given volume of stellar fluid is determined by the interaction between neutrino opacities and the composition, the later of which is dependent on the choice of equation of state. The neutrino–matter interactions include emission, absorption, and non-isoenergetic scattering on free nucleons (Reddy et al. 1998) with weak magnetism corrections (Horowitz 2002); emission/ absorption (electron capture) on nuclei

([Langanke et al. 2003](#)); isoenergetic scattering on nuclei, including ion–ion correlations; non-isoenergetic scattering on electrons and positrons; and pair emission from e^+e^- annihilation ([Bruenn 1985](#)) and nucleon–nucleon bremsstrahlung ([Hannestad and Raffelt 1998](#)).

4.2.4 Nuclear Physics Inputs

In order that the hydrodynamics, nuclear transmutations, and neutrino transport be tied closely to the thermodynamics, the equation of state (EoS) must be invoked every cycle, in fact, a number of times each cycle. Furthermore, the EoS must provide not only the quantities needed for the hydrodynamics, e.g., pressure, internal energy and entropy as a function of density, temperature, and electron fraction, but the element composition, and chemical potentials as well, as these latter are needed for the neutrino transport. WeakLib was first developed to use the CHIMERA canonical LS EOS in tabular form. This was in order to verify that simulations using direct calls to the LS EOS and using tabular data resulted in the same outcome, with no numerical errors arising from the use of the table. The WeakLib infrastructure replaced direct queries to the EOS, as well as interpolation and EOS inversion queries.

4.2.5 Simulation Parameters

All simulations were made with a polar grid with 720 adaptive radial zones and 240 angular zones for the 2D runs. The angular zones are of fixed angle, but the radial zones change in size in order to meet several stability criteria, namely the Courant-Friedrichs-Lewy condition, and to sufficiently resolve the shock at the PNS surface. During the evolution, the radial zones are gradually and automatically repositioned during the remap step to follow changes in the radial structure. For each angular zone, the 720 radial zones from the center to the edge of the grid form an radial “ray”. This particular domain decomposition is dictated by the nature of the neutrino transport, which is radial in its current implementation. The boundary

conditions at each pole are reflecting; motion transverse to pole is cancelled, whereas radial motion is not. This leads to elongation of structures at the poles; this is an issue only in 2D, not in 1D or 3D. Some inner zones are frozen during collapse (due to necessity of Courant conditions).

In order to explore the impact of the EOS in CCSNe, we ran both 1D and 2D simulations for each EOS, with the CHIMERA default LS/BCK EOS serving as our fiducial case. Additionally, we forked the SFHo EOS simulation run into two runs, one with the electron fraction capped at the upper limit of the SFHo table ($Y_e = 0.6$) and another for which the dependent EOS quantities had been extrapolated up to $Y_e = 0.7$ to investigate the effect this would have on the simulation.

For the purpose of this study, we chose to take the initial conditions for each simulation from a pre-supernova $15 M_\odot$ progenitor of [Woosley and Heger \(2007\)](#), in order to best investigate the collapse of a massive star in the central region of the mass range associated with core-collapse supernovae. This initial conditions of this progenitor are those of a star evolved from the main sequence to the onset of core collapse. The simulations all begin with the iron core of the progenitor of a mass of $1.34644 M_\odot$ at the point of collapse.

Chapter 5

EOS in CCSN Simulation

The choice of equation of state has significant implications/consequences for the supernova dynamics. As the star collapses, the formation and evolution of the proto-neutron star (PNS) is characterized by the equation of state in several ways. The stiffness of a particular equation of state (as well as the rate of deleptonization) influences the mass profile and compactness of the forming PNS, and it furthermore governs the initial character of the PNS through core bounce. Additionally, the equation of state impacts the formation and subsequent evolution of the supernova shock profile. The equation of state also influences the radial position of the neutrinosphere, the location of which can significantly impact the spectra of the emitted neutrinos. As the equation of state determines the conditions by which the nuclear composition evolves, so it will likewise affect the conditions for nucleosynthesis.

While many aspects of core-collapse supernovae have been studied extensively in simulations, I will focus here on discussion of simulations and studies performed to investigate the sensitivity of supernova simulations to the equation of state. Colgate and White first proposed the neutrinos might be what powers core-collapse supernovae explosions (Colgate and White 1966). They were also among the first to produce numerical simulations of such events. Initially, most studies of the impact of the EOS in supernova were limited to investigations in spherical symmetry (one-dimensional models). 1D models generally failed to

produce an explosion due to a lack of sufficient energy transfer between the neutrinosphere and the stalled shock layers in the delayed reheating mechanism. This failure to explode colors the conclusions of such studies. The effects of incompressibility, collapse, and the neutrino signature were studied by both [Thompson et al. \(2003\)](#) and [Suwa et al. \(2006\)](#), [Suwa et al. \(2007\)](#), the former comparing different incompressibility versions of the LS EOS ($K = 180, 220, \text{ and } 375$), and the latter comparing the same versions of the LS EOS to the STOS EOS. [Hempel et al. \(2012\)](#) studied the sensitivity of spherically symmetric simulations to the EOS via their new group of RMF EOSs (FSUGold, IUFSU), in terms of both the hydrodynamics as well as neutrino signals. [Steiner et al. \(2013\)](#) found that central densities at bounce show a weak correlation with the symmetry energy slope parameter L and that there is a moderate influence of the symmetry energy on the evolution of the electron fraction in the PNS, and introduced the new SFHx and SFHo EOSs into comparison. EOS studies of multi-dimensional supernova models have shown that the gravitational wave signal is highly sensitive to the EOS ([Kotake et al. 2003, 2012](#); [Scheidegger et al. 2010](#); [Richers et al. 2017](#)). Variations in the neutrino luminosity due to EOS differences were also investigated, and it was also found that the choice of EOS influence the time of onset of explosion ([Marek and Janka 2009](#)). Investigations in two dimensions by [Suwa et al. \(2013\)](#) compared the STOS and LS EOSs and found that the softer LS EOS leads to a more energetic explosion.

5.1 CHIMERA Results

In this section I will discuss the main differences in hydrodynamics and neutrino emission properties between the simulations using different EOSs, and then focus in the following sections on the specific epochs that define the evolution of the supernova event. While the EOSs utilized in these simulations are constrained to the best degree currently possible, they nonetheless still contain significant uncertainties as indicated by their variations. The goal of this section is to observe and understand how the differences in the EOS models affect the outcomes and observables in these simulations.

The effect of multidimensionality cannot be understated in terms of our understanding of the CCSN explosion mechanism; multidimensionality allows continued accretion and explosion to coexist, something impossible in one-dimensional simulations. Effects of the EOS on PNS formation and bounce, as well as the early and late accretion phases only become clear via multidimensional simulations. Firstly, post-shock non-radial flow allows for convection, which enhances neutrino heating compared to one dimensional simulations. Additionally, convective buoyant plumes and the SASI, neither possible in one dimension, expand the shock, increasing the residence time of the matter in the gain region and lengthening the advection timescale. All these multidimensional features further invigorate the neutrino-driven shock expansion. The coexistence of accretion and outflow also allows for mixing, which allows the chemical composition of the matter to evolve. The difference in the shock radius evolution between the one- and two-dimensional models (see Figures 5.1 and 5.2) shows that for all EOSs used in this dissertation, dimensionality dictates the success or failure of the explosion. For the one-dimensional models, the shock radii expand, with the LSBCK model's shock radius expanding the furthest to 176 km before stalling just before 100 ms after bounce, at which point the shock radius of the LSBCK model and all other EOS models begin to contract and explosion is never achieved, as shown in Figure 5.1. The one and two-dimensional models evolve similarly until shortly after 100 ms after bounce, when the shock expansion for the two-dimensional models is revived, and the shock radii continue to increase monotonically for all two-dimensional models over the duration of the simulations, indicating successful explosions, as shown in Figure 5.2.

One frequently used parameter to gauge the evolution of the shock expansion is the time required for the mean shock radius of a given model to reach 500 km. Generally by this time, the shock is lifting the SFHx model achieves a mean shock radius of 500 km first at 251 ms after bounce, followed by the LSBCK model at 261 ms, then by DD2 at 265 ms, SFHo at 270 ms, FSUGold at 279 ms, IUFSU at 280 ms, and lastly by NL3 at 284 ms, which is ~ 30 ms later for the same WH07 progenitor than was required for the LSBCK model of

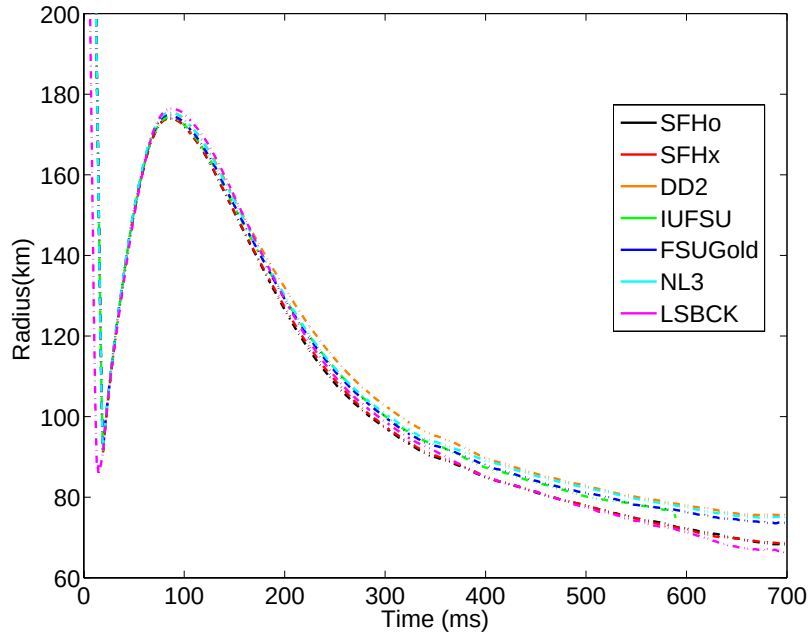


Figure 5.1: 1D Shock Radius vs. Time. The mean shock radii for the 1D models can be seen rising initially and then peaking near 100ms and then receding for all models.

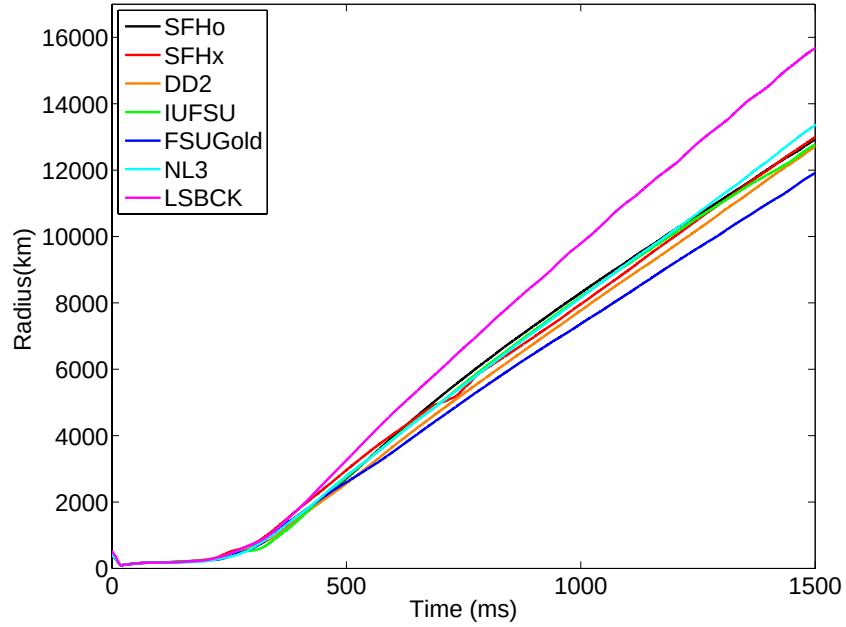


Figure 5.2: 2D Shock Radius vs. Time. The mean shock radii for the 2D models can be seen rising monotonically for all models.

the earlier B-series CHIMERA simulation (Bruenn et al. 2013, 2016), indicating the range of impact of other changes.

One indicative quantity frequently used to gauge the evolution of the explosion energy is the diagnostic energy, E^+ , the integral by volume of the total energy density, the combination of the kinetic (e_{kin}), thermal (e_{th}), and gravitational (e_{grav}) energy densities:

$$e_{\text{tot}} \equiv e_{\text{kin}} + e_{\text{th}} + e_{\text{grav}} \quad (5.1)$$

over the zones between the shock and PNS for which the total energy density is positive $e_{\text{tot}} > 0$ (Buras et al. 2006; Suwa et al. 2010; Müller and Janka 2014; Bruenn et al. 2013, 2016). In this definition, the thermal energy is defined as the internal energy density with the rest mass energy of conserved particles subtracted out. The diagnostic energy becomes positive for all EOS models around 200 ms and differentiates for each model over the course of the simulation, as shown in Figure 5.3. This specific measure of explosion energy is useful as it is reflective of the total energy transferred into the central engine of the neutrino-reheating mechanism. The total energy evolves in a similar fashion for each model, but does not become positive until 387 ms after bounce for the LSBCK model, the first model to become energy positive, followed by SFHo at 451 ms after bounce, SFHx at 478 ms, DD2 at 491 ms, NL3 at 526 ms, IUFSU at 540 ms, and lastly by FSUGold at 584 ms, as shown in Figure 5.4.

The choice of equation of state also has significant effects on the evolution of the neutrino luminosities. All models maintain a similar electron neutrino luminosity profile during the few ms after bounce, climbing and peaking during the neutrino burst at shock breakout, and decline rapidly until about 100 ms after bounce. After this point, the effects of the differences in infalling matter accretion rates and the structure of each model’s proto-neutron star affect the luminosities of all neutrino species, influencing the post-breakout peak and subsequent rate of decline in the luminosity of all neutrino species, as shown in Figures 5.5 through 5.8. The luminosity of all models can be seen declining from the neutrino burst at shock

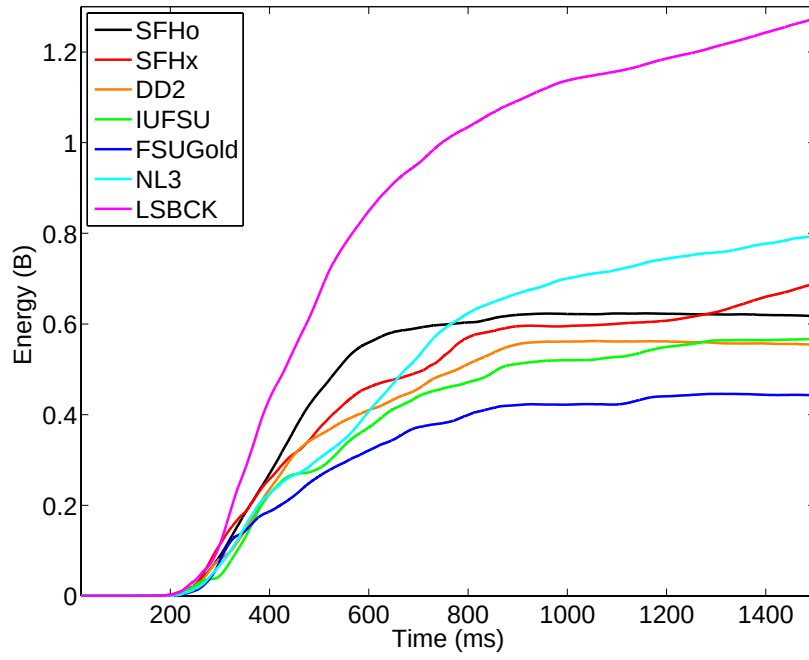


Figure 5.3: Diagnostic Energy vs. Time

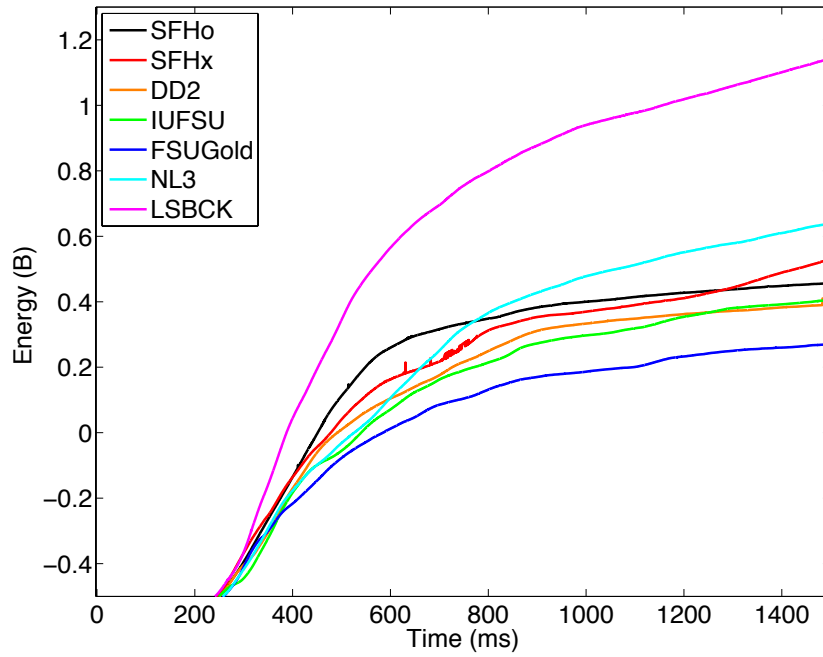


Figure 5.4: Total Energy vs. Time

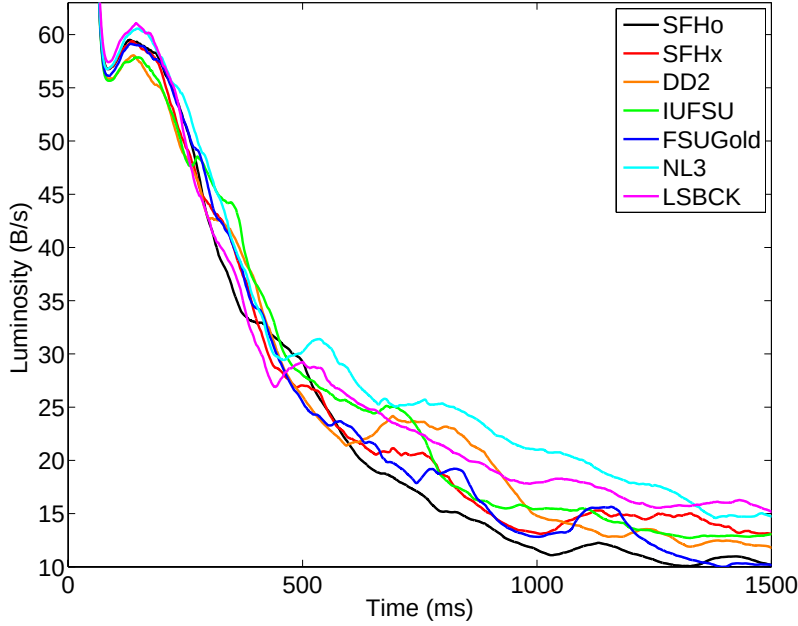


Figure 5.5: Electron Neutrino Luminosity vs. Time.

breakout, then rebounding during the mantle accretion phase. Around 400 ms, the rate of change in luminosity decreases for all neutrino species due to the shift to emission from the PNS core.

After 300 ms differences in mass accretion rate onto the PNS lead to divergence in the proto-neutron star masses for the individual models, as seen in Figure 5.9. The accretion at the gain radius is shown in Figure 5.10. The LSBCK model, having been more luminous in the 100 - 200 ms period, experiences less gain accretion after that period as a result, though accretion returns after 500 ms. The models with more late gain accretion tend to develop larger neutron stars. For the same mass, a stiffer EOS leads to a larger PNS radii, as seen in Figure 5.11, with lower central density, as shown in Figure 5.12. The EOSs that result in the most compact (PNS mass divided by radius) PNSs are (up until 1000 ms) the LSBCK model, followed by the SFHo and SFHx models, which become the most compact after 1000 ms, as seen in Figure 5.13

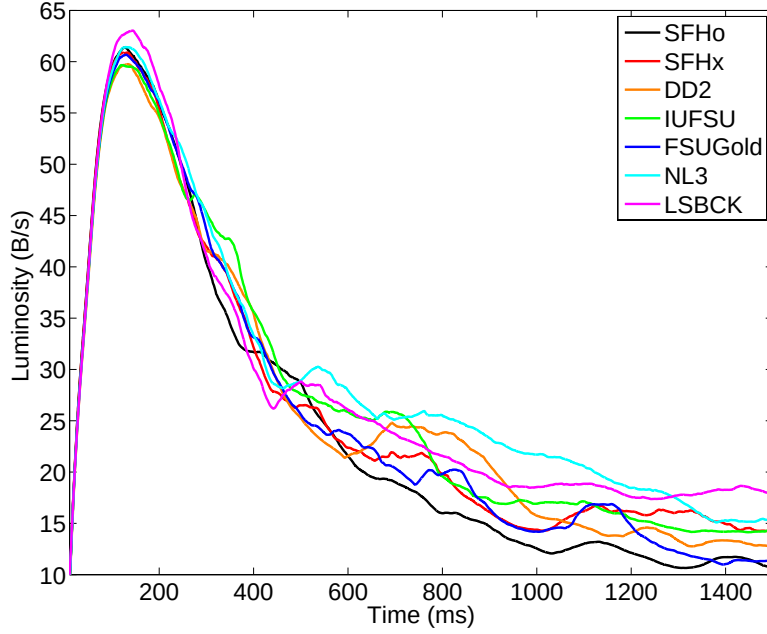


Figure 5.6: Electron Antineutrino Luminosity vs. Time. The magnitude of the luminosities for electron antineutrinos in each model are slightly higher than for electron neutrinos beyond the shock burst at breakout.

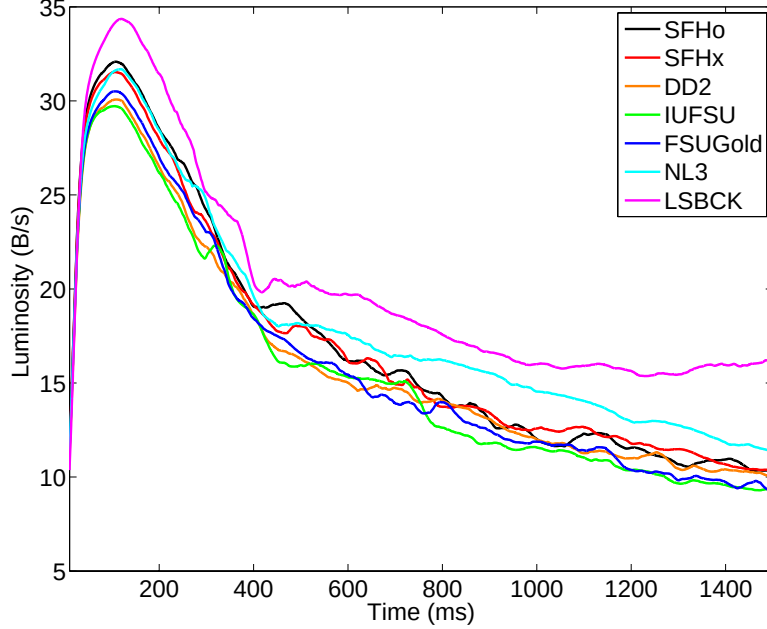


Figure 5.7: Mu and Tau Neutrino Luminosity vs. Time. The LSBCK model exhibits a higher relative peak luminosity for the μ and τ neutrino than the other models. Also, the post-400ms period also shows relatively higher luminosity with respect to the other models.

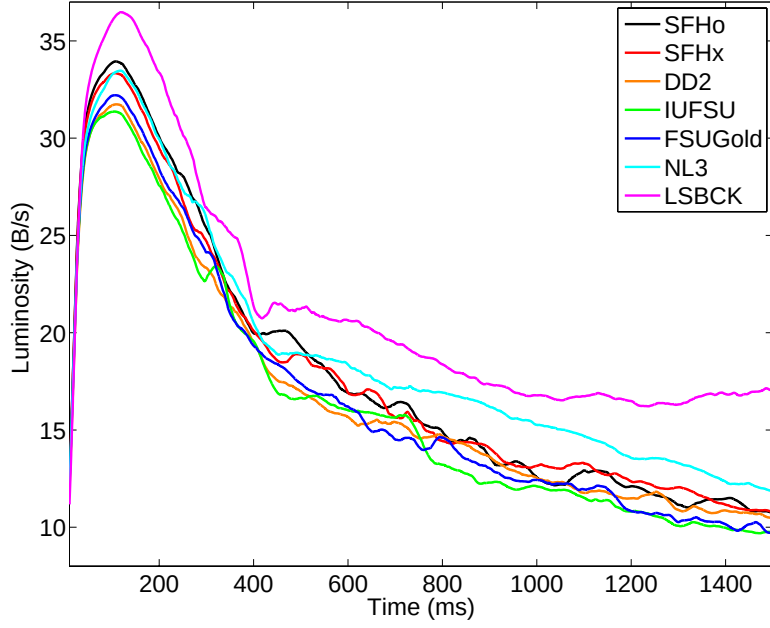


Figure 5.8: Mu and Tau Antineutrino Luminosity vs. Time. Again, the magnitude of the luminosities for μ and τ antineutrinos in each model is slightly higher than for μ and τ neutrinos, and the relative luminosities for LSBCK are again higher than for the ν_e and $\bar{\nu}_e$.

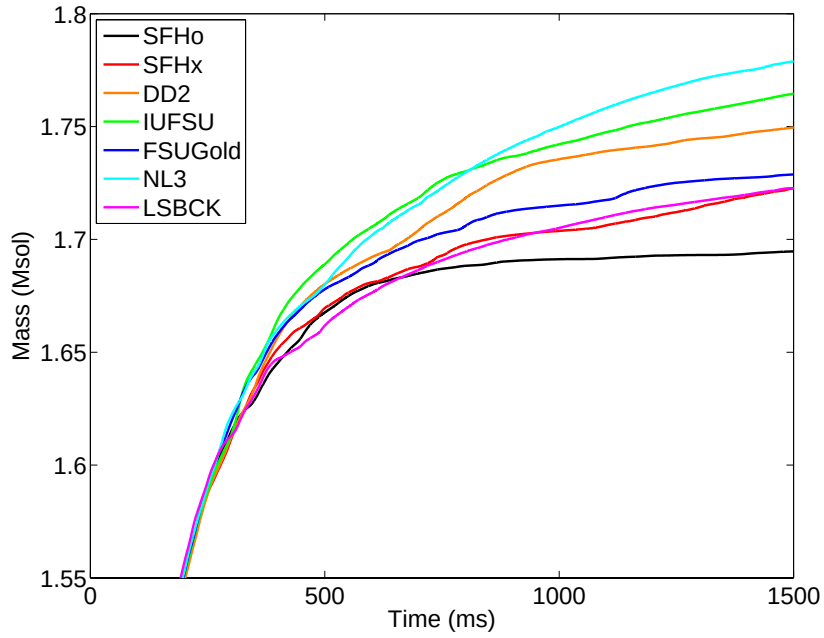


Figure 5.9: PNS Mass vs. Time.

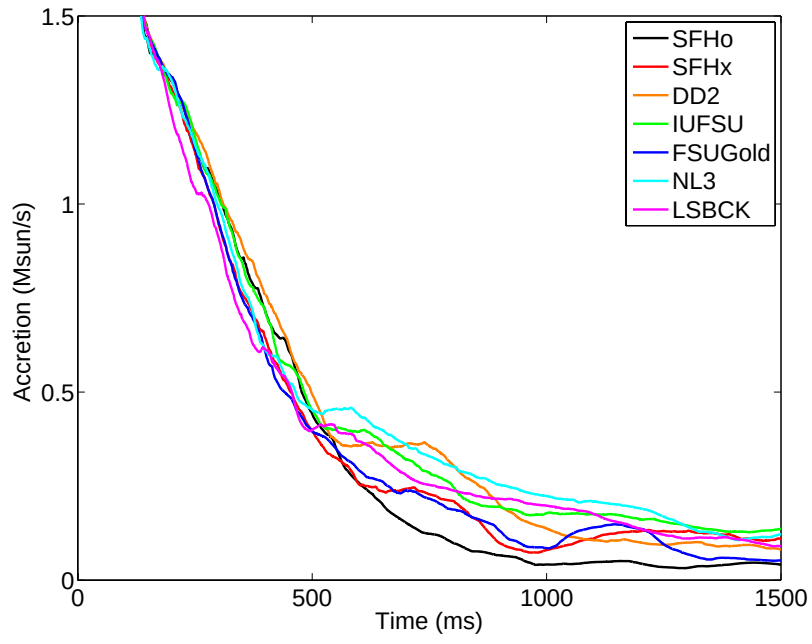


Figure 5.10: Gain Accretion vs. Time.

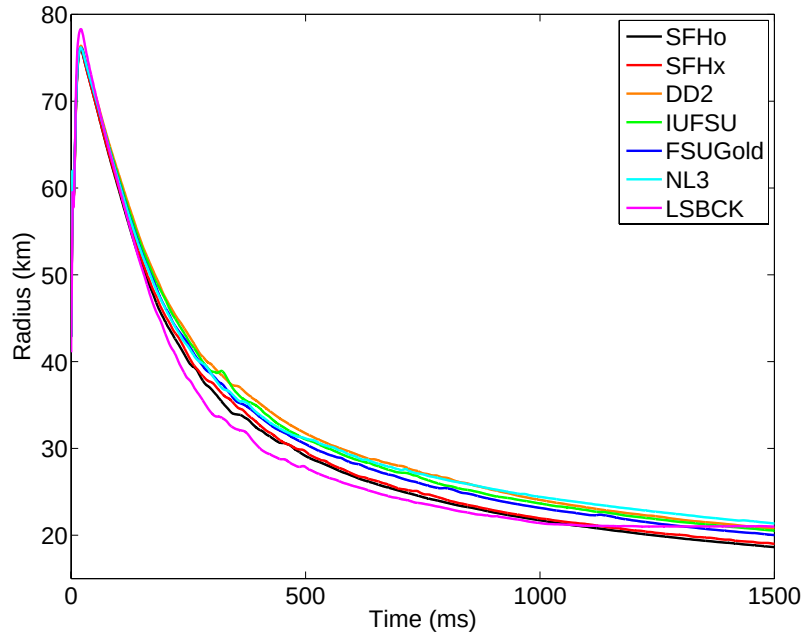


Figure 5.11: PNS Radius vs. Time.

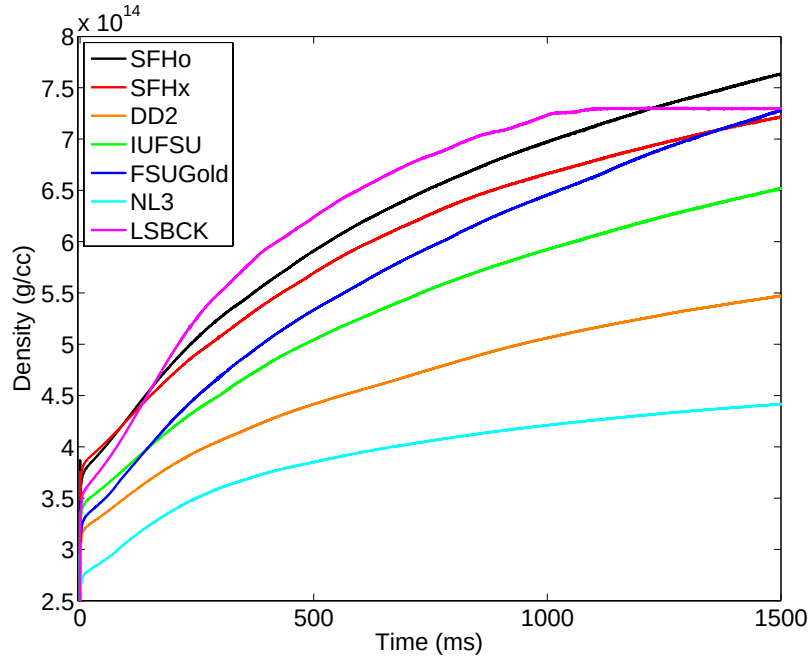


Figure 5.12: PNS Central Density vs. Time.

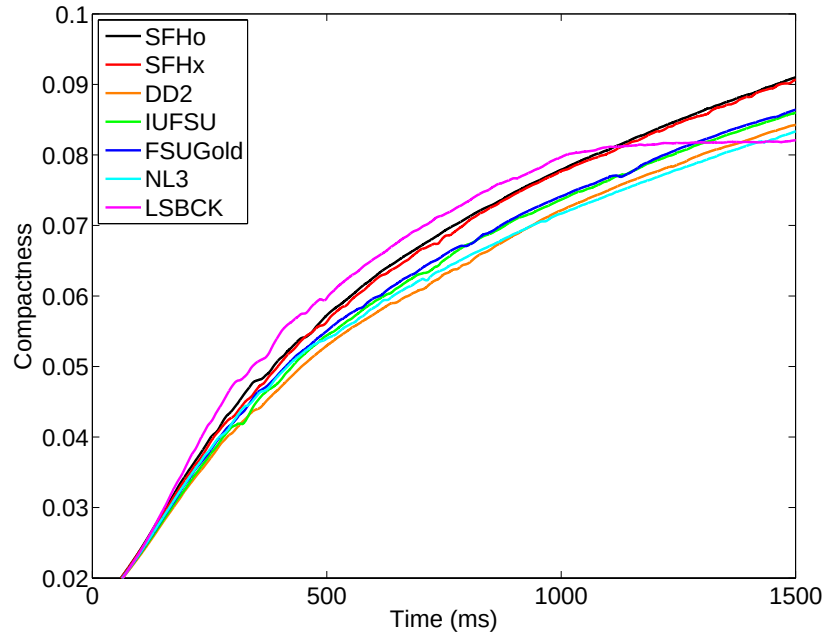


Figure 5.13: PNS Compactness vs. Time. The evolution of the PNS compactness, the ratio of the enclosed mass to the radius.

5.2 Effects of EOS during Collapse

As the star collapses, electron captures on both free protons and on protons within nuclei, the unique composition given by each EOS in the collapse, determine the deleptonization and subsequent neutrino emission rates. Through this sensitivity, the dynamics of collapse and bounce are highly dependent on electron and neutrino capture on nuclei, especially nuclei with masses $A > 60$ (Hix et al. 2004). During collapse, the dominate pressure comes from the degenerate electron gas, and electron captures further reduce this. As collapse continues, increasing densities are reached in the inner core and the core deleptonization continues until at eventually densities of about $1 \times 10^{11} \text{ g cm}^{-3}$ the emitted neutrinos become trapped, as their free-path length at these densities becomes less than the size of the core. This leads to the slowing of deleptonization as re-absorption of the neutrinos on neutrons begins to approach emission rates, and the beta processes balance each other. Bounce begins as the core approaches a density of $1 \times 10^{14} \text{ g cm}^{-3}$. The deleptonization history of the core during collapse affects initial PNS mass as the sound speed and hence the sonic point depend on the electron density, with LSBCK having the lowest initial mass and radius as bounce begins, as seen in Figure 5.14 and Figure 5.15.

5.3 Effects of EOS at Bounce

Approaching bounce, the central core material stiffens by nuclear repulsion as nuclear density is approached and exceeded, and the compressed inner core's collapse is halted and reversed, with the dominant pressure contribution arising from the nucleons and nuclei. At this point, a bounce shock is formed at the sonic point that separates the homologous inner core and the infalling outer core into which it rebounds. The evolution of our models during collapse is for the most part self-similar (with the same general character independent of the collapse time), with LS bouncing slightly later (1.2 ms) than the RMF models. Differences appear at bounce, with the NL3 EOS has the shallowest density profile, while the SFHx and SFHo

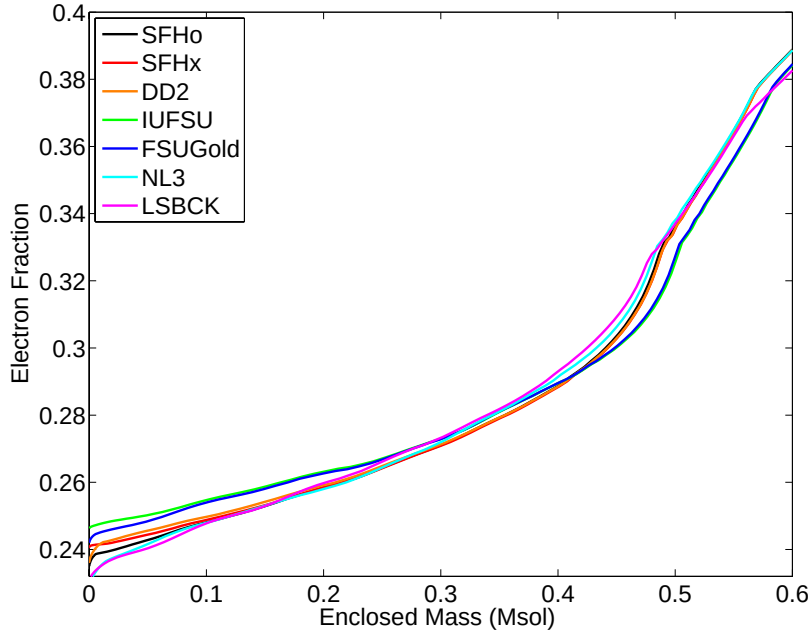


Figure 5.14: Electron Fraction vs. Enclosed Mass, at $\rho_{\text{central}} = 1 \times 10^{14}$

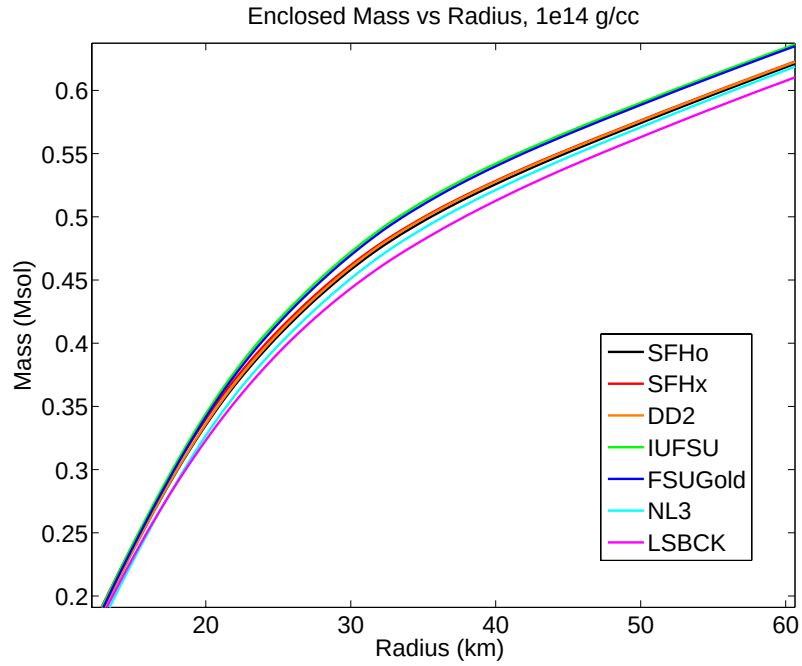


Figure 5.15: Enclosed Mass vs Radius, at $\rho_{\text{central}} = 1 \times 10^{14}$

Table 5.1: Comparison of EOS Effects at Bounce

At Bounce	SFHo	SFHx	DD2	IUFSU	FSUGold	NL3	LSBCK
Homologous Inner Core (Msol)	0.491	0.478	0.485	0.493	0.496	0.519	0.490
Central Density ($\times 10^{14}$ g cm $^{-3}$)	3.873	3.521	3.114	3.279	3.450	2.728	3.351
Central Electron Fraction	0.243	0.238	0.244	0.247	0.251	0.250	0.245
Central Entropy Per Baryon	1.054	1.054	1.053	1.041	1.044	1.062	1.13
PNS Surface Electron Fraction	0.295	0.289	0.293	0.292	0.292	0.305	0.298
PNS Surface Entropy Per Baryon	2.98	2.51	2.50	2.55	2.65	3.50	3.26
Minimum Initial Radius (km)	46.18	46.19	46.83	46.11	46.10	45.99	44.65

models have the steepest, as seen in Figure 5.16. The NL3 has the lowest pressure at the surface of the homologous core, as seen in Figure 5.17. The PNS in the LSBCK model contracts more with a lower initial radius than the other models, $\sim 5\%$ smaller than for the RMF EOS models. As shown in Figure 5.18 the entropy has a maximum value for each model at the coincident PNS surface and shock location, with the NL3 model having the highest shock entropy $3.50 k_b/\text{baryon}$, followed by the LSBCK model $3.26 k_b/\text{baryon}$, then the remaining models. The entropy value descends behind this surface, as the core material is unshocked, with a notable differentiation at $\sim 0.7 M_\odot$ corresponding to the transition from the nuclear LS EOS to the lower-density BCK EOS at $\rho = 1 \times 10^{11}$ g cm $^{-3}$. The LSBCK model has a notably hotter PNS surface than the RMF models, but a cooler interior, as seen in Figure 5.19. The infall velocity of material above the inner core approaches $\sim 0.3 c$, as seen in Figure 5.21, which emphasizes the location of the homologous inner core boundary, as does the electron fraction (Figure 5.22). The NL3 model has the largest inner core mass at bounce, in part due to its higher incompressibility. Figure 5.23 shows the forming neutrinospheres within the core for each model as well as the deleptonization still occurring in the infalling material pre-shock up to the higher density region of neutrino trapping where a region of low electron neutrino luminosity forms.

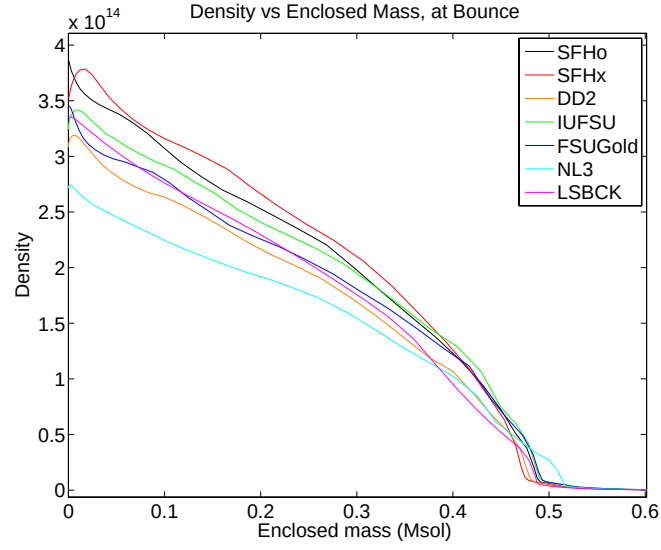


Figure 5.16: Density (g cm^{-3}) vs. Mass

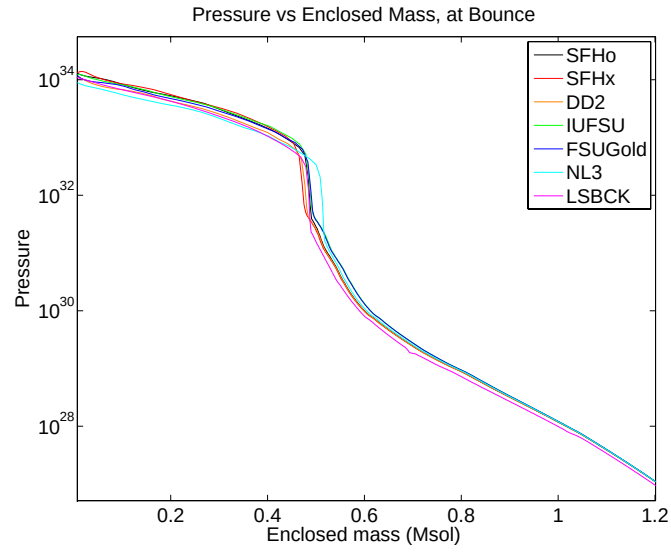


Figure 5.17: Pressure (dynes/cm^2) vs. Enclosed Mass

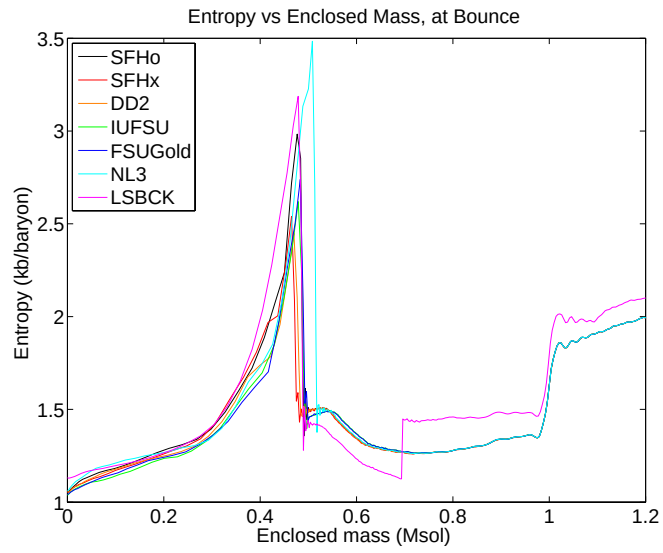


Figure 5.18: Entropy vs. Enclosed Mass

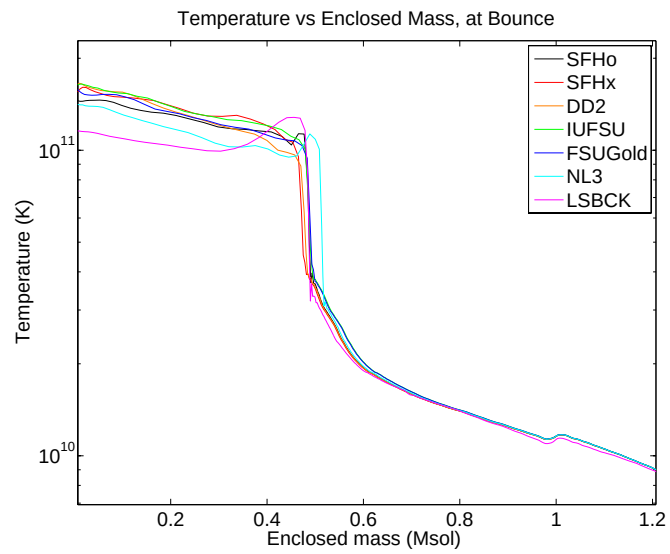


Figure 5.19: Temperature vs. Enclosed Mass

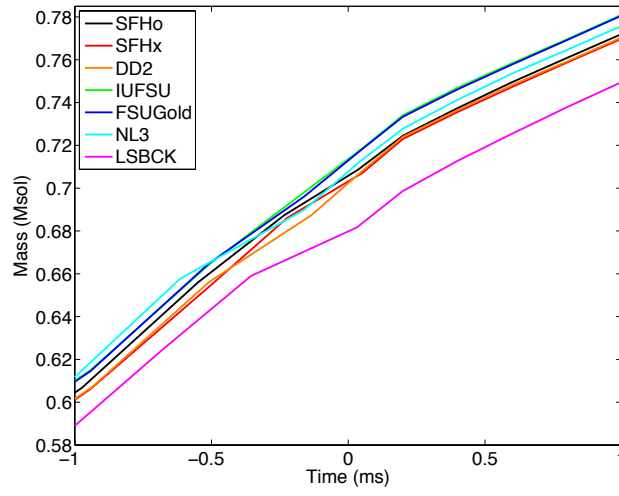


Figure 5.20: PNS Mass vs. Time, at Bounce

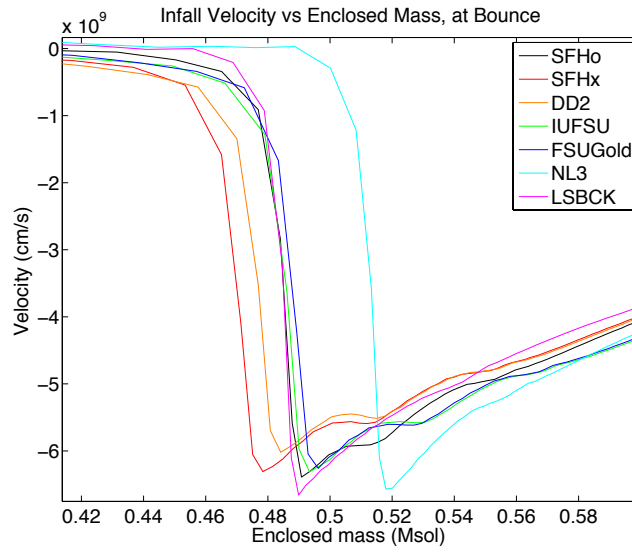


Figure 5.21: Infall velocity vs. Mass

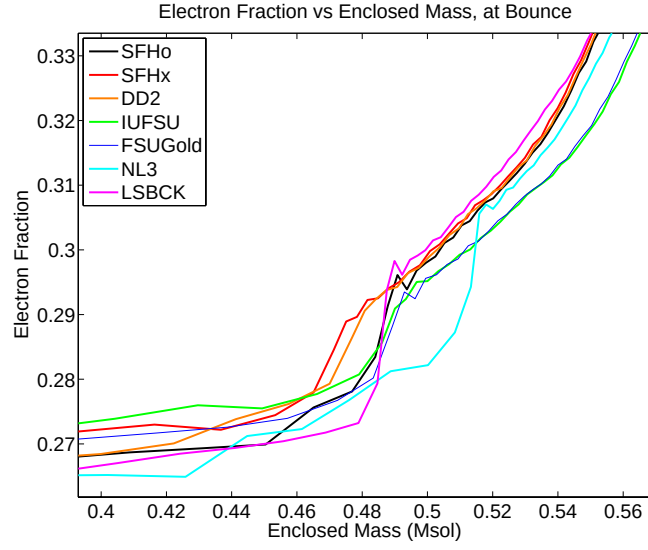


Figure 5.22: Electron Fraction vs. Mass

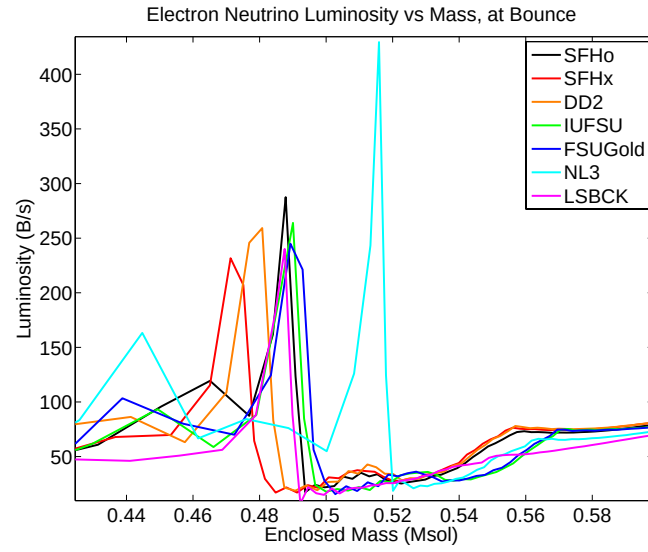


Figure 5.23: Electron Neutrino Luminosity vs. Mass, at Bounce

5.4 Initial Accretion Phase

After bounce, the bounce shock proceeds outward into the infalling material. The different conditions and history of the inner core lead to different initial shock locations and shock evolution for each model. Even before bounce occurs, the rate of ν_e emission begins rising, as the increasing high densities reached during the beginning of bounce may provide increased electron capture efficiency. As shown in Figure 5.24, the luminosity dips briefly just after bounce ($\sim 4 - 5$ ms), as the enhanced compression from the rising shock and infalling matter cause a temporary condition of increased opacity. The vast amount of trapped ν_e produced from trapping to this point are contained in the dense post-shock matter, until the rising shock reaches densities low enough that the diffusion speed of the ν_e is faster than the shock propagation speed. When the shock reaches this neutrino-transparent matter, a luminous burst of electron neutrinos is released, with the peak occurring at $8.4 - 8.5$ ms after bounce for the RMF EOS models and at 8.8 ms for the LSBCK model, with the LSBCK model having $\sim 4\%$ less luminosity than the RMF group on the rising edge and top of the peak, as seen in Figure 5.24. In general, the stiffness of the equation of state correlates with the height of the luminosity peak, highlighted in Figure 5.25, with the softer equations of state trapping more neutrinos and hence releasing fewer at shock breakout. Concurrently with the rising peak at shock break out, charged-current processes in the shock-heated mantle of the PNS lead to the production of ν_e and $\bar{\nu}_e$ as well as heavy-lepton neutrinos, $\nu_{\mu,\tau}$ and $\nu_{\mu,\tau}^-$. The μ and τ neutrinos have the highest root mean square (RMS) energies, followed initially by the ν_e then $\bar{\nu}_e$, though the $\bar{\nu}_e$ soon surpass the ν_e , as seen in Figure 5.26. The ν_e initially have higher RMS energy than the $\bar{\nu}_e$ and $\nu_{\mu,\tau}$, but after bounce, as the shock heats the infalling matter which heats the mantle of the PNS, the RMS energy of the $\bar{\nu}_e$ and $\nu_{\mu,\tau}$ (generated by pair production) rise past that of the $\bar{\nu}_e$.

Figure 5.27 documents the rapid expansion of the shocks in all models through. The softer nuclear EOS of the LSBCK model has lead to a smaller initial shock and PNS radius at bounce. For the first few milliseconds after bounce, the shock of LSBCK model expands

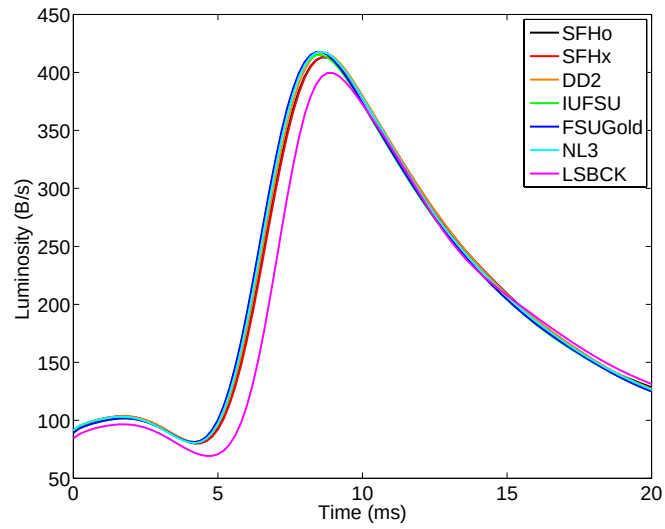


Figure 5.24: Electron Neutrino Luminosity vs Time

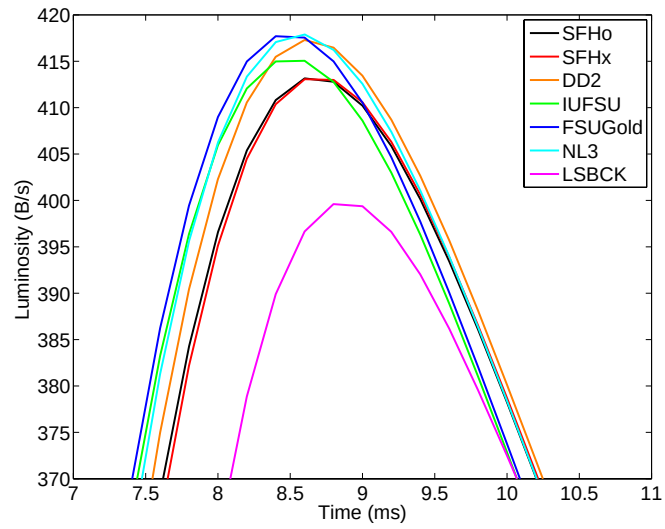


Figure 5.25: Electron Neutrino Luminosity Peak vs Time, zoomed view

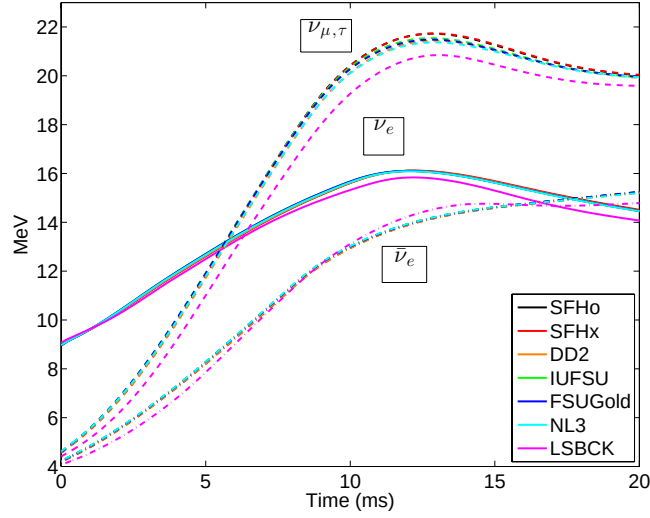


Figure 5.26: Neutrino RMS Energy vs Time. The ν_e are the solid line, the $\bar{\nu}_e$ the dash-dot line, and $\nu_{\mu,\tau}$ are the dashed line.

more slowly than the other models. However the energy stored in this compression leads to a steeper rebound in radius, and by ~ 7 ms after bounce, the LSBCK model’s shock radius exceeds that of the rest. This effect is mirrored in the temporarily expanded PNS radius, as seen in Figure 5.28, with maximum PNS radius being reached around 20 ms after bounce in all models. Despite its distinctly larger PNS radius during this period, the LSBCK model’s PNS mass lags that of the other models during this early epoch, as can be see in Figure 5.29.

For the first 50 ms after bounce, the shock continues to propagate outward and the electron neutrino luminosity diminishes from the high value at the breakout burst and levels off at $\sim 50 \text{ B s}^{-1}$ (Figure 5.30). During this later period of electron neutrino luminosity decline ($t_{pb} < 30$ ms), the LSBCK model is slightly more luminous in this aspect than the other models (Figure 5.30). The electron antineutrino luminosities continue to increase monotonically, with SFHo and SFHx slighter higher, and LSBCK slightly lower, as seen in Figure 5.31.

During this phase in the simulation, material is accreting through the shock and is significantly shock-heated, and a hot “bubble” has formed behind the shock. The “gain

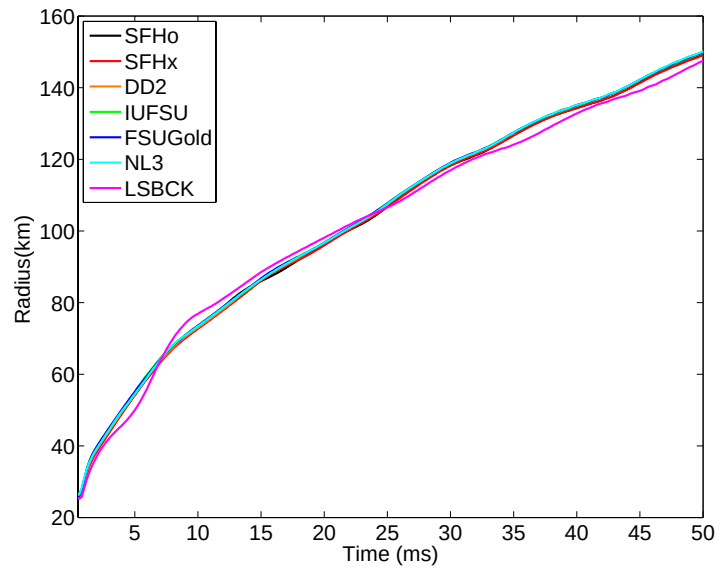


Figure 5.27: Shock vs Time

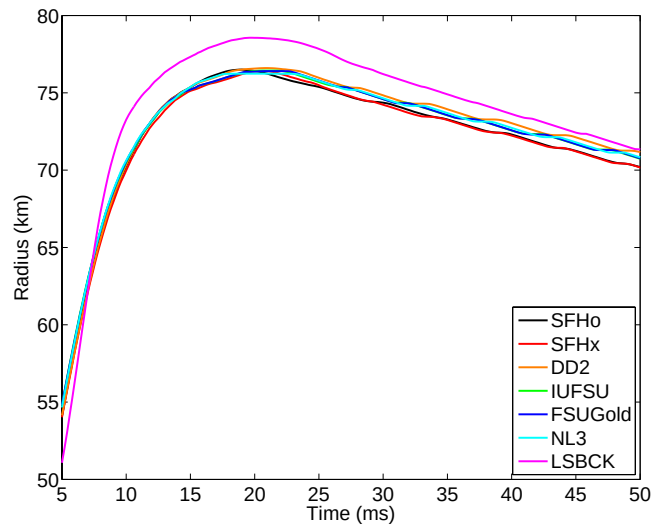


Figure 5.28: Proto-neutron Star Radius vs Time

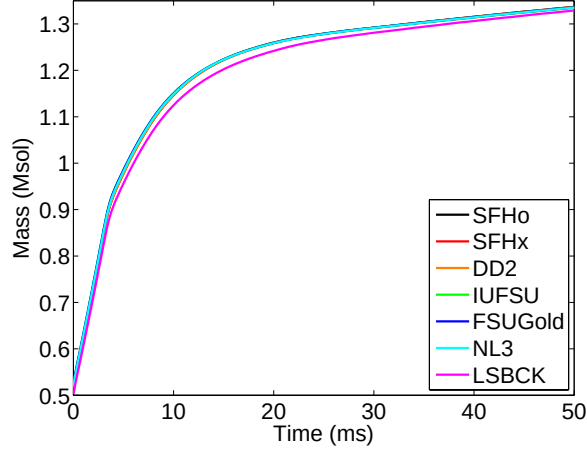


Figure 5.29: Proto-neutron Star Mass vs Time

radius” in a simulation, the radius going from inside the PNS outward to the edge of the star at which the heating becomes more significant than the cooling is a noteworthy landmark.

In this early epoch (before 50 ms), there exists a region of neutrino cooling behind the shock but above a region of neutrino heating, and a second very thin gain region just behind the shock, as seen at 30 ms in the left panel of Figure 5.32. Here we consider only four of the models, for visual clarity, but all RMF models behave similarly to SFHo, NL3, and FSUGold. The LSBCK model experiences greater neutrino cooling in this central post-shock region due to a higher percentage of alpha particles, and correspondingly smaller abundance of free nucleons (see right panel of Figure 5.32). As the alpha particles have lower neutrino interaction cross sections than the sum of their individual nucleons (due to the alpha particle being “double magic”, i.e. absorption is energetically disfavored), when they represent a larger percentage of the fluid by mass, the fluid becomes more neutrino transparent and heating becomes less efficient. At 40 ms (see Figure 5.33) this continues, but the cooling is lessened. By 50 ms, the outer cooling region has transitioned into a lesser heating region, as shown in Figure 5.34, and the demarcated gain radius, specifically detected as the point where the net neutrino heating becomes negative going from the outside inward, now shows the expected gain radius. This produces the discontinuity shown in Figure 5.35. The discontinuity occurs 3 ms later for LSBCK which continues to exhibit a slightly larger

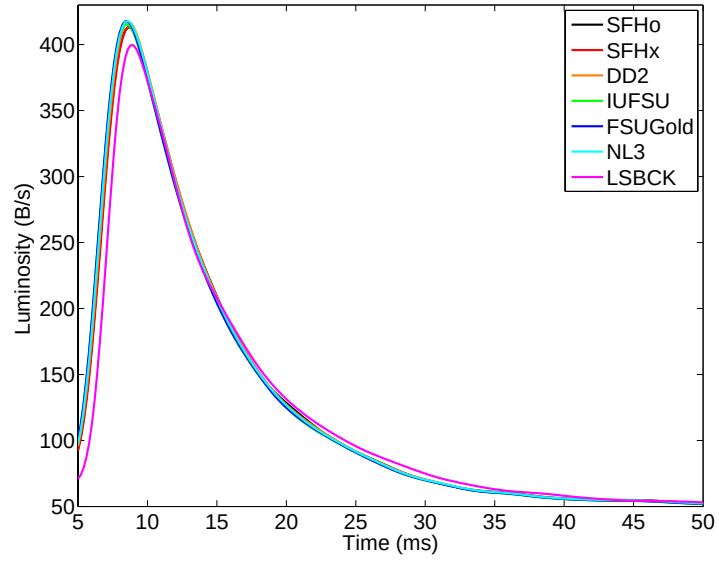


Figure 5.30: Early Electron Neutrino Luminosity vs Time

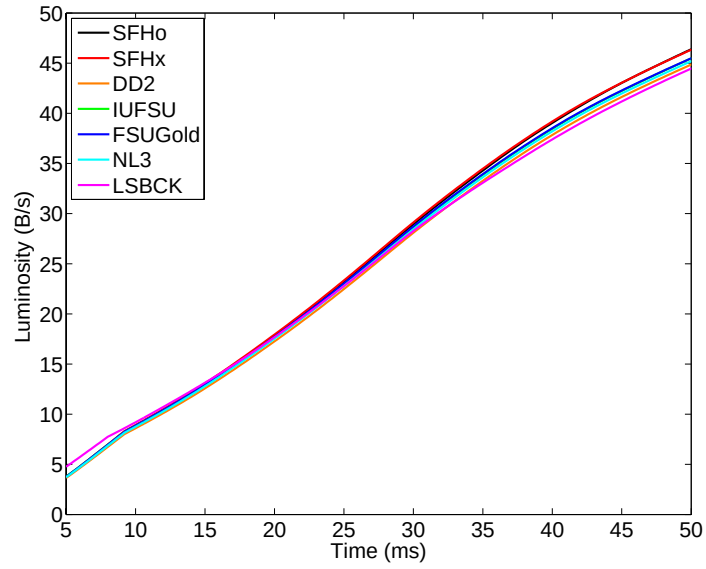


Figure 5.31: Early Electron Antineutrino Luminosity vs Time

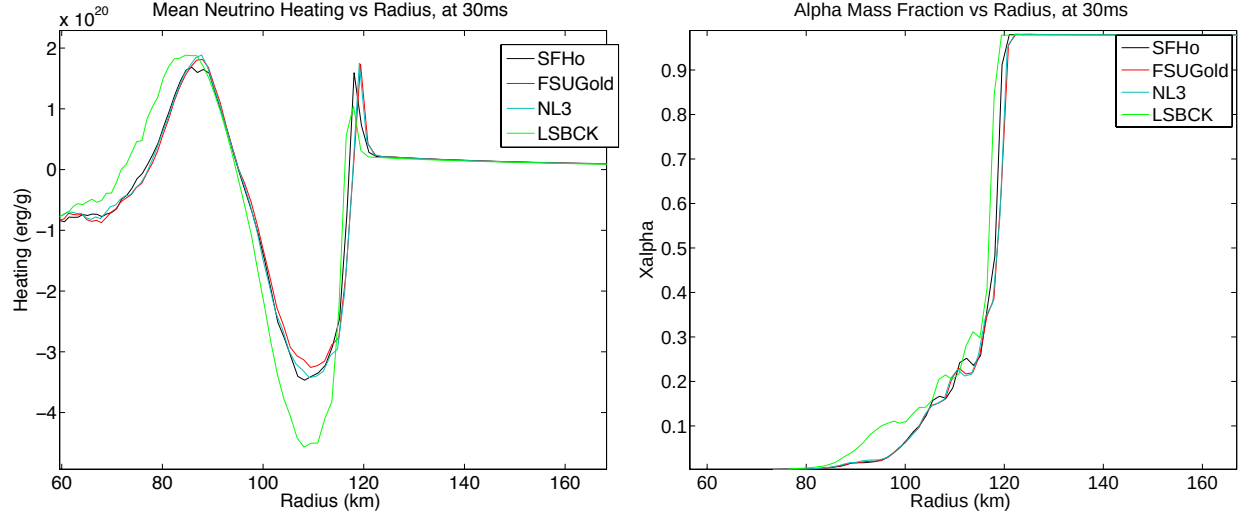


Figure 5.32: Alpha Mass Fraction and Neutrino Heating vs. Radius, 30ms.

α fraction and slightly reduced heating in this region. From that point forward, the two gain regions have joined, and the conventional inner neutrino cooling region/ outer neutrino heating region paradigm is established.

5.5 Onset of Convection

After the shock stalls, a negative entropy gradient is formed in the post-shock layer, as neutrinos heat the accreted material in the gain region. For the 2D models, neutrino-heated matter in the inner region of the gain region rises outward in buoyant plumes, pushing aside colder infalling material which can in turn fall toward the region of stronger neutrino heating to be heated. Images of the entropy 50 ms after bounce illustrate a time when neutrino-driven convection is just beginning. Convection develops somewhat earlier in the SFHo model (Figure 5.36), with the onset of convection in the SFHx (Figure 5.37), DD2 (Figure 5.38), IUFSU (Figure 5.39), and FSUGold (Figure 5.40) models following, and lastly in the fiducial LS-BCK model (Figure 5.42). This convective development is occurring despite the similarity in the neutrino luminosity at this time (see Figure 5.30 and 5.31, indicating that this is possibly an effect of the EOS in the relatively low density heating

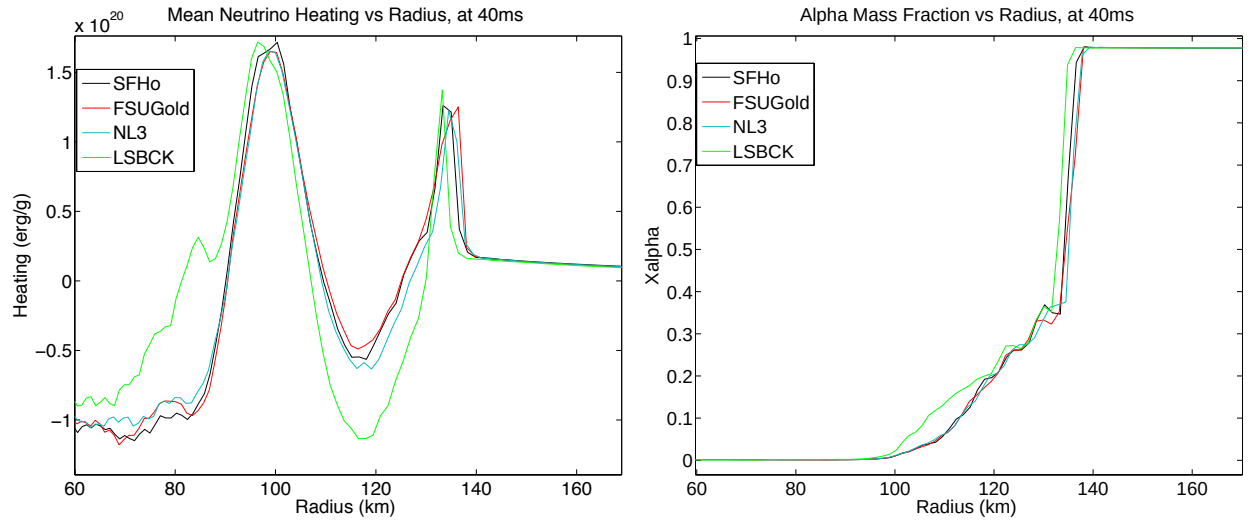


Figure 5.33: Alpha Mass Fraction and Neutrino Heating vs. Radius, 40 ms.

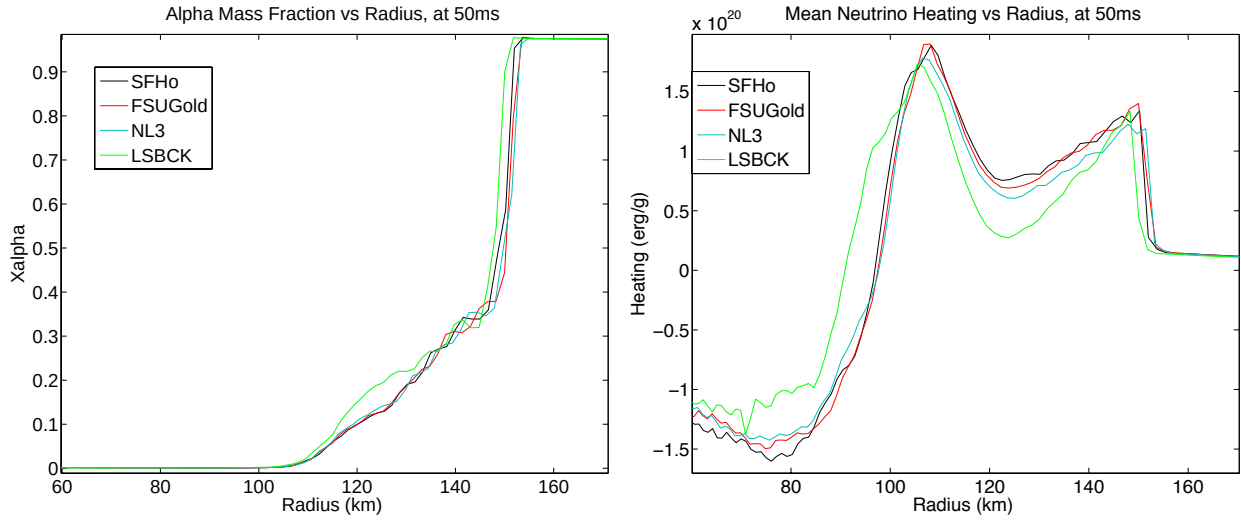


Figure 5.34: Alpha Mass Fraction and Neutrino heating vs. Radius, 50 ms.

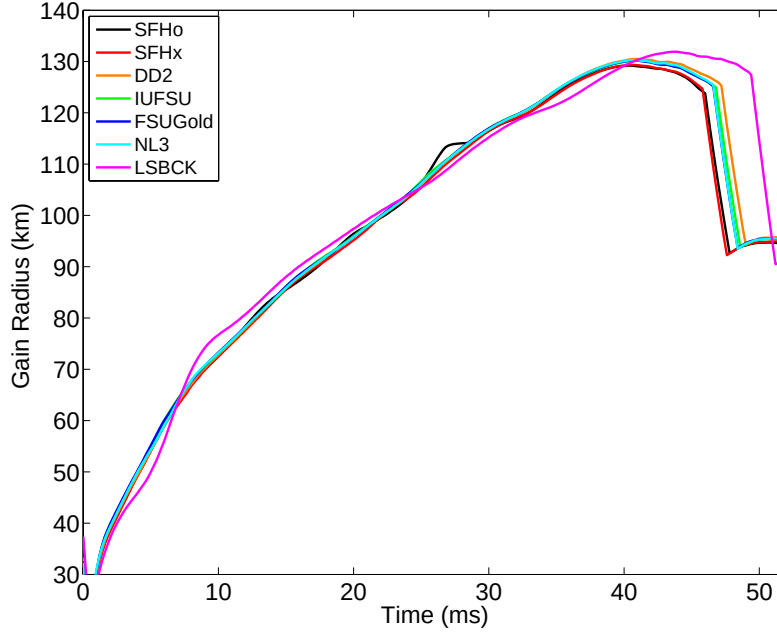


Figure 5.35: Gain Radius vs Time, Post-bounce

region, rather than a difference originating deep in the nuclear matter in the PNS. Previous studies have shown that larger mass accretion rates through the gain surface inhibit the onset of convection (Müller and Janka 2014; Bruenn et al. 2016). As shown in Figure 5.43, the SFHo model does indeed have the lowest accretion rate in the early quiescent accretion period, while the LSBCK model has the highest, with the other models in between, though more like SFHo. Convection in the gain region corresponds with a steepening in the rise in the gain volume, as shown in Figure 5.44. A larger gain region lengthens the timescale for material to advect through the gain region.

The shock radius continues to expand for all models (see Figure 5.45) monotonically until ~ 65 ms, at which point the rise begins to falter. Supersonically infalling matter is compressed as it comes through the shock, depositing its former gravitational potential energy, turned into kinetic energy, into heat behind the shock. While the the shock wave has been expanding in each model it has been losing energy as the high temperature behind the shock leads to the disassociation of heavy nuclei into nucleons and α particles as they pass through the shock front. Combined with the energy lost to the initial neutrino burst

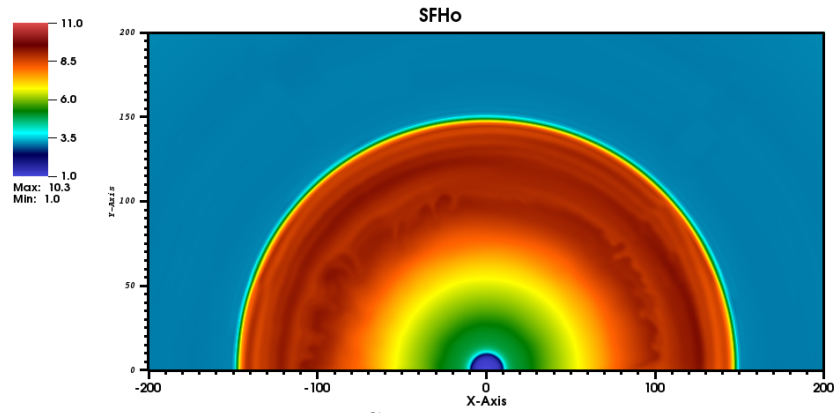


Figure 5.36: SFHo Entropy, 50 ms

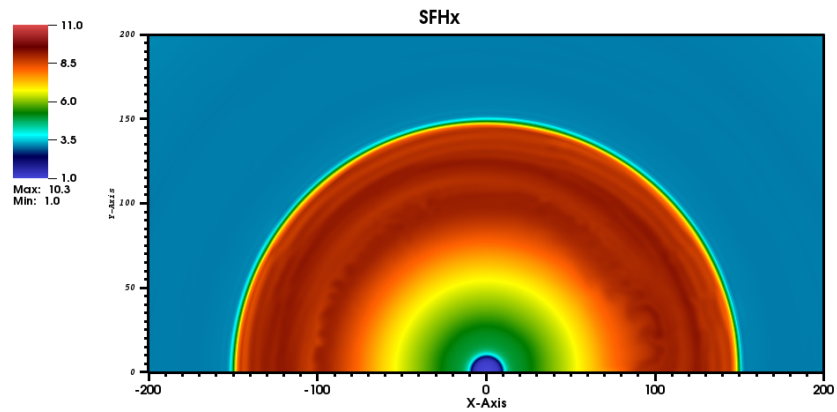


Figure 5.37: SFHx Convection, 50 ms

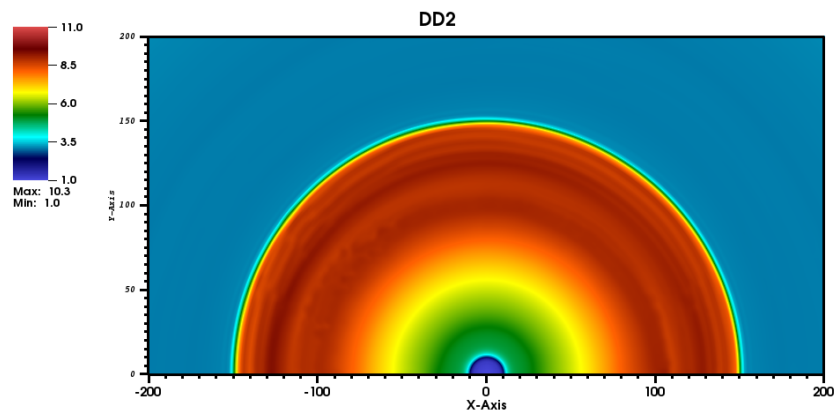


Figure 5.38: DD2 Entropy, 50 ms

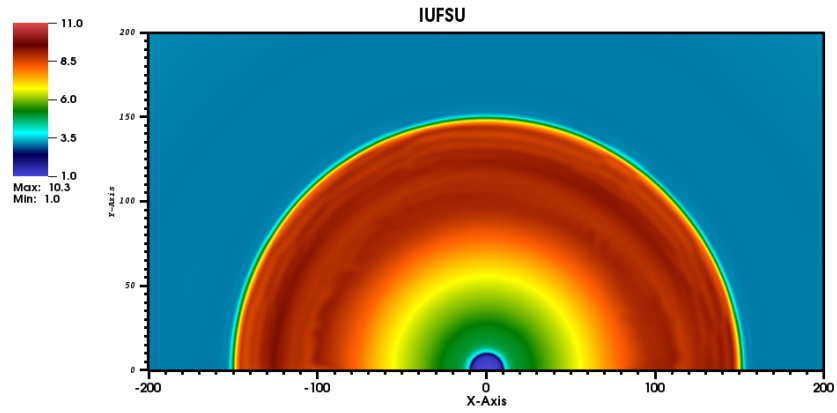


Figure 5.39: IUFSU Entropy, 50 ms

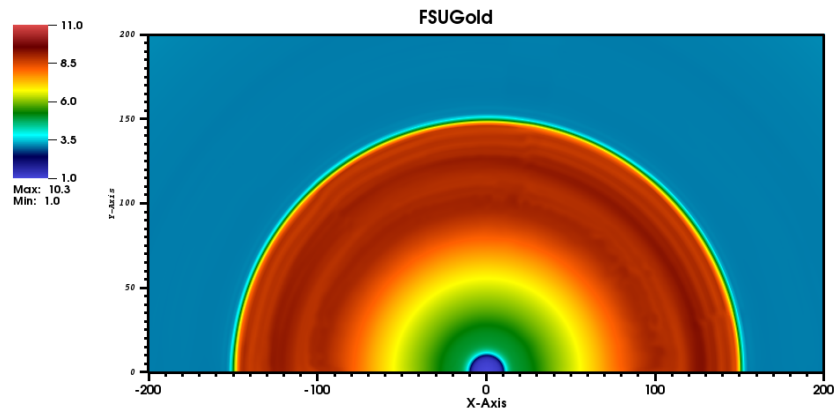


Figure 5.40: FSUGold Entropy, 50 ms

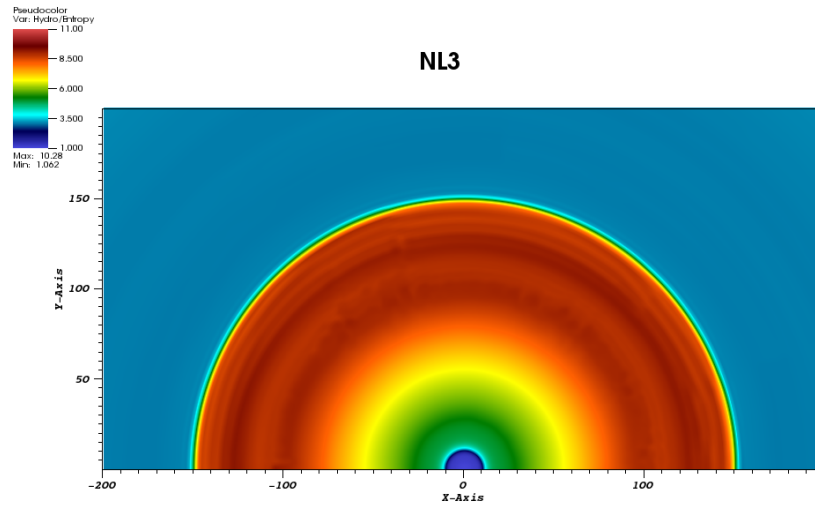


Figure 5.41: NL3 Entropy, 50 ms

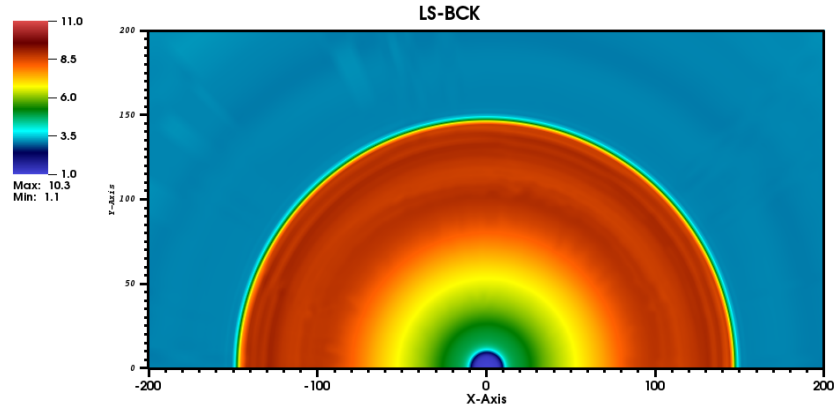


Figure 5.42: LSBCK Entropy, 50 ms

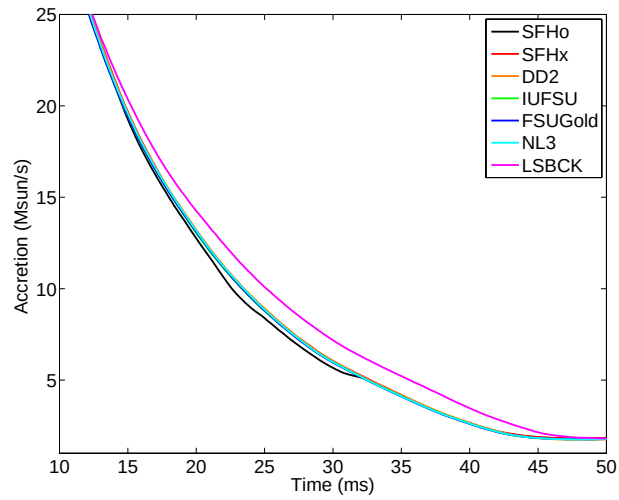


Figure 5.43: Gain Radius Mass Accretion vs. Time, 50 ms

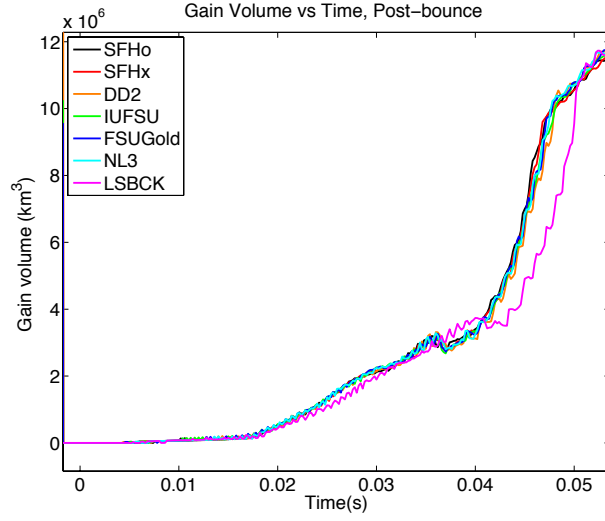


Figure 5.44: Gain Volume vs. Time, 50 ms

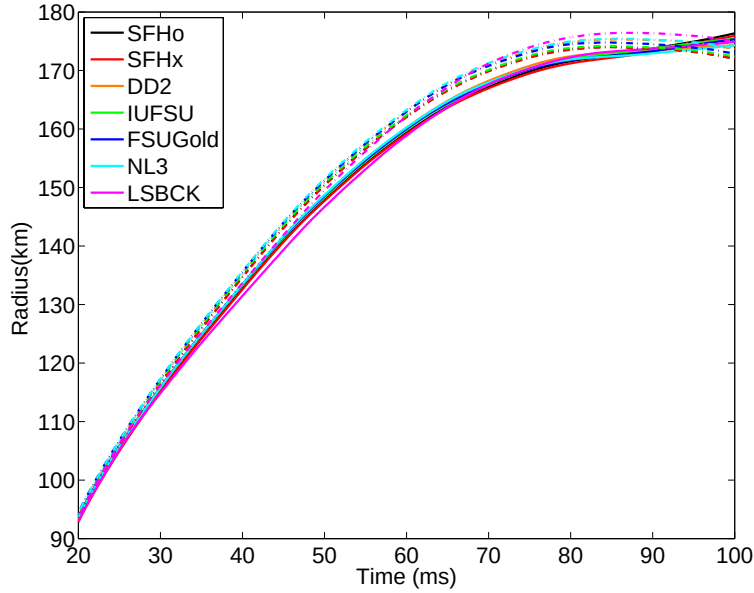


Figure 5.45: Shock Radius vs Time, 1D vs 2D. The one-dimensional shock trajectories can be seen here as the dotted-dashed lines while the solid lines represent the two-dimensional model shock trajectories. The 1D model shock radii peak between 80–90 ms, and start to decline, while the 2D model shock radii are beginning to rise.

at shock breakout, this energy drain causes the shock to stagnate and become a stationary accretion shock, just above a radius of 170 km in all of these models. At this point it is clear that the “prompt mechanism” for supernova explosion is not sufficient to explode the star in the models for any EOS tested here.

5.6 Effects of EOS by 100ms

In the period of time between 50 and 100 ms after bounce, the accretion through the shock and resulting deceleration has been increasing the amount of matter in the now convecting gain region, which is being heating by energy from the ν_e and $\bar{\nu}_e$ fluxes as well as heavy-lepton neutrinos, $\nu_{\mu,\tau}$ and $\nu_{\mu,\tau}^-$ radiated by the PNS. The neutrino interactions depositing the bulk of the emitted energy in the material in the gain region are

$$\nu_e + n \rightarrow p^+ + e^- \quad (5.2)$$

$$\bar{\nu}_e + p \rightarrow n + e^+ \quad (5.3)$$

By 100 ms, robust, neutrino-driven convective overturn is underway, and can be seen in all models (Figure 5.46), and while all models are exhibiting turbulent convection, the beginning of asphericity can be seen in the upper left and middle left panels, where rising hot plumes have impacted through the shock front. It is around this time in the simulation that the 1D and 2D shock locations begin to diverge, as seen in Figure 5.45. For the 1D models, convection cannot occur, and hence cannot aid the heating mechanism. For the 2D models, neutrino heating in the gain layer has caused the temperature and pressure behind the shock to increase to the point that dynamic ram pressure on the shock from the rising turbulent convective plumes causes the revival of the stalled shock to begin. At 100 ms after bounce we see the first instances of the structures of the turbulent convection (specifically rising, buoyant Rayleigh-Taylor plumes expanding into the shock) causing the development of shock asphericity, most notably in the SFHo model (see upper left panel of Figure 5.46).

The appearance of asphericity causes an increase in advection timescale in the aspherical regions, aiding in heating. The RMS energies for all neutrino species, initially lower for LSBCK, have converged and rise monotonically (see Figure 5.47).

The combination of the expanded width of the gain region (Figure 5.48) along with the rising ν_e and $\bar{\nu}_e$ luminosities (Figure 5.49 and Figure 5.50) lead to a rise in the heating efficiency η_{heat} (Figure 5.51), the ratio of the neutrino energy deposition rate in the gain layer $\dot{Q}_\nu$ divided by the sum of the ν_e and $\bar{\nu}_e$ luminosities at the gain surface $L_{\nu,\text{gain}}$. The η_{heat} of the LSBCK model can be seen to exceed that of the RMF models for virtually all of the time between bounce and 100 ms after bounce. An additional indication of the shock revival can be found in looking at the advective and heating behavior of each model. A useful parameter to characterize the advection behavior is the advection timescale, τ_{adv} , which specifically denotes the dwell time of a parcel of matter in the gain region. In the time between 50–100 ms, the mass residing in the gain region begins to rapidly increase, which increases the advection timescale and heating efficiency. In concert with τ_{adv} , description of the heating of such accreting matter can be characterized by defining a heating timescale, τ_{heat} :

$$\tau_{\text{heat}} = \frac{E_{\text{th}}}{\dot{Q}_\nu} \quad (5.4)$$

where E_{th} is the thermal energy of the material contained in gain region. Past studies have shown that when the ratio of these timescales, $\tau_{\text{adv}}/\tau_{\text{heat}}$ approaches and exceeds unity, shock revival is imminent (Thompson et al. 2005; Buras et al. 2006; Bruenn et al. 2016). As can be seen in Figure 5.52, the $\tau_{\text{adv}}/\tau_{\text{heat}}$ ratio is increasing monotonically for all models, at rates that will achieve unity soon after 100 ms. The shock radius experiences an upward inflection just before 100 ms for all models, as seen in Figure 5.45. The properties of the PNS at 100 ms can be seen in Table 5.2.

Table 5.2: Comparison of EOS Effects at 100 ms

At 100 ms	SFHo	SFHx	DD2	IUFSU	FSUGold	NL3	LSBCK
PNS Mass (Msol)	1.430	1.429	1.425	1.428	1.427	1.427	1.428
Central Density ($\times 10^{14}$ g cm $^{-3}$)	4.269	4.257	3.515	3.805	3.750	3.072	4.159
Central Electron Fraction	0.243	0.238	0.246	0.250	0.253	0.257	0.253
Central Entropy Per Baryon	1.054	1.054	1.053	1.041	1.044	1.062	1.13

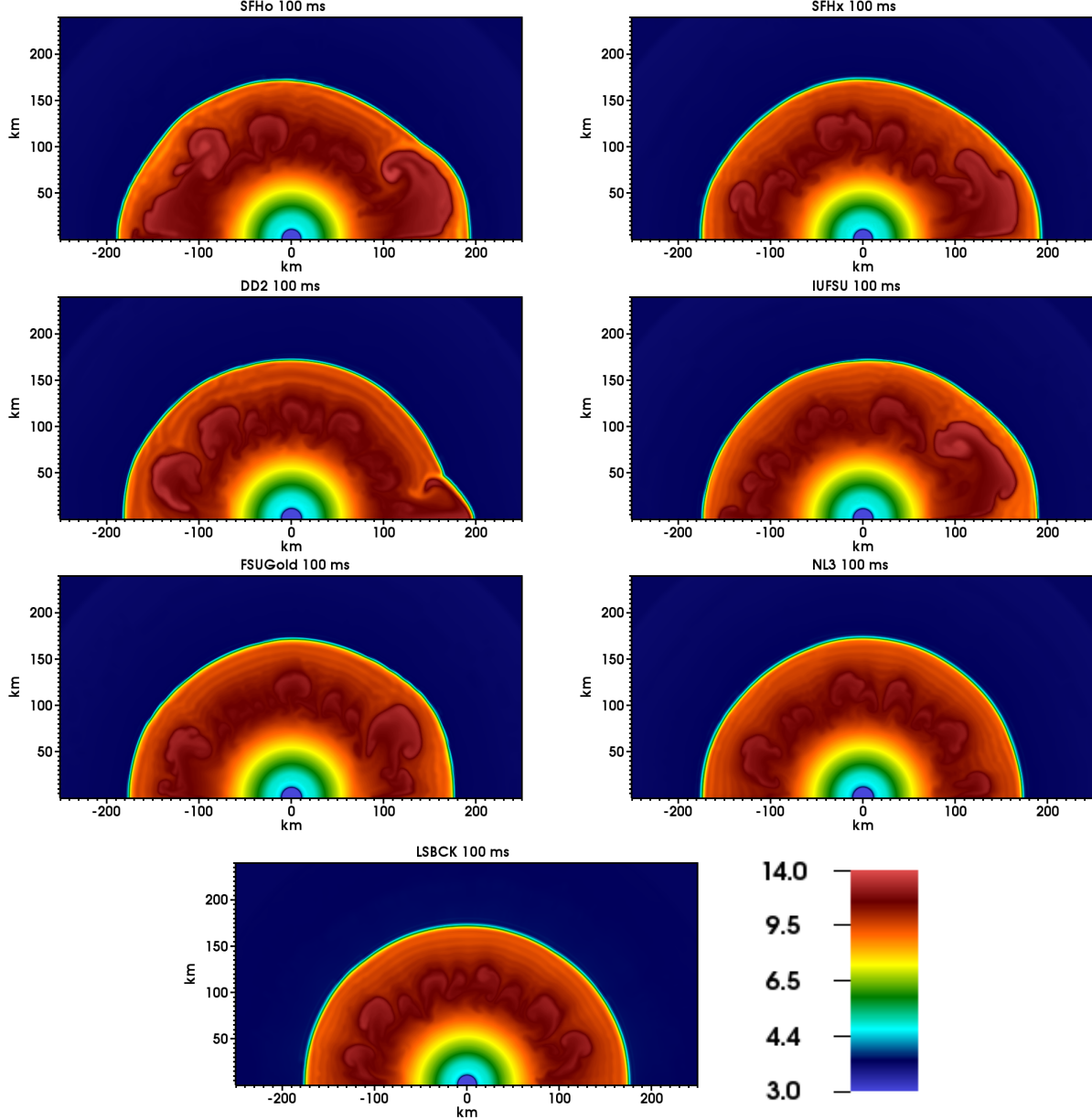


Figure 5.46: Entropy (log scale), 100 ms

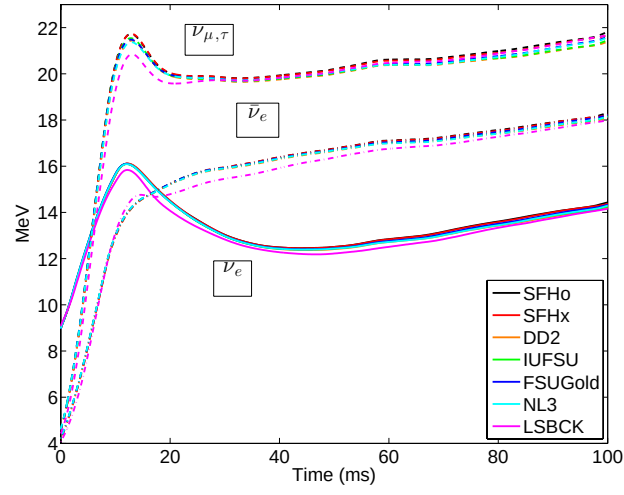


Figure 5.47: Neutrino RMS Energy vs Time

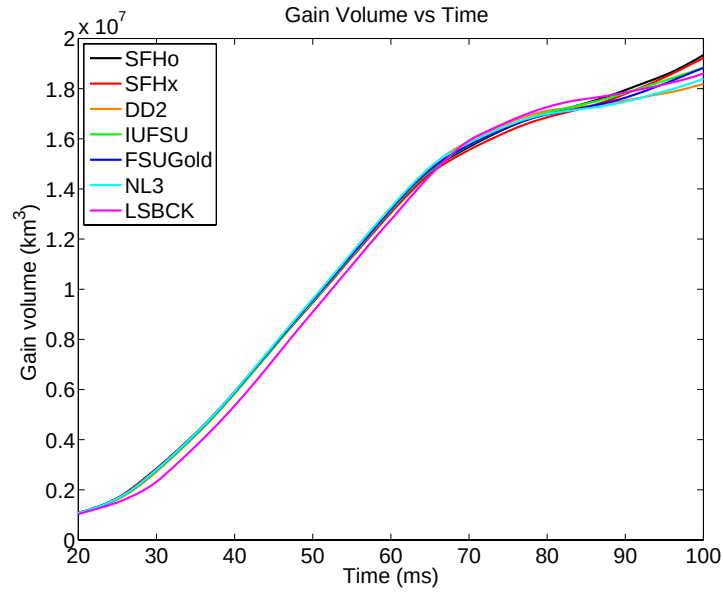


Figure 5.48: Gain Volume vs. Time, 100 ms

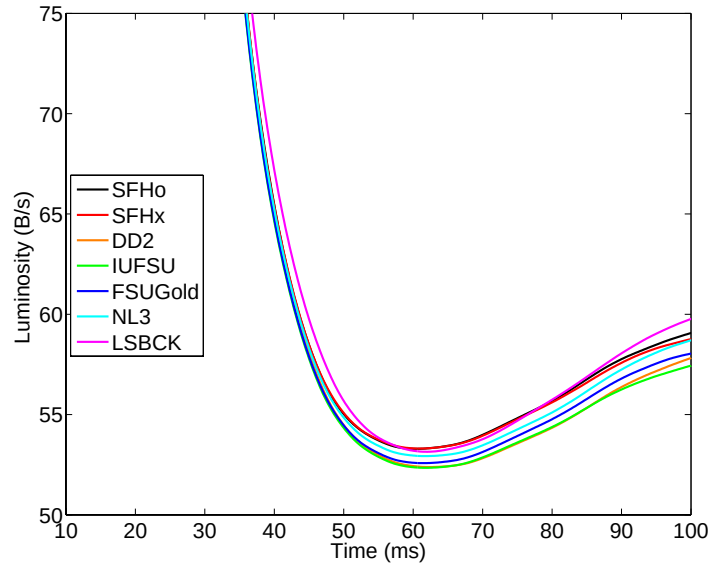


Figure 5.49: Electron Neutrino Luminosity vs Time, 100 ms

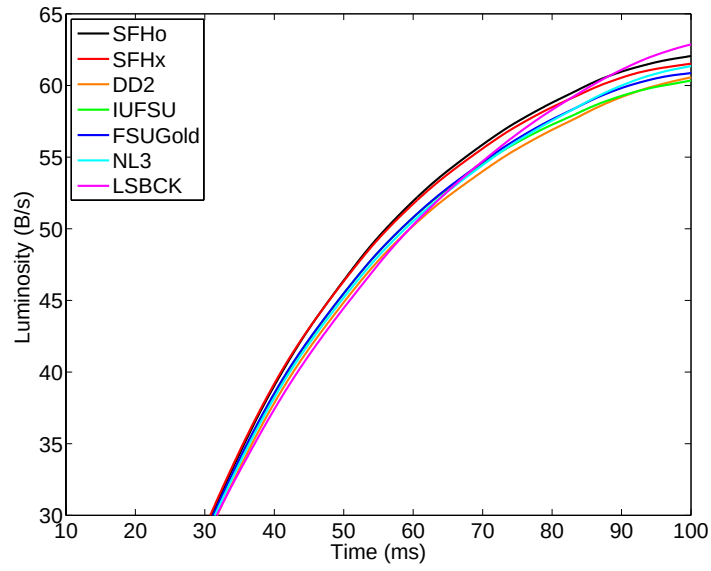


Figure 5.50: Electron Antineutrino Luminosity vs Time, 100 ms

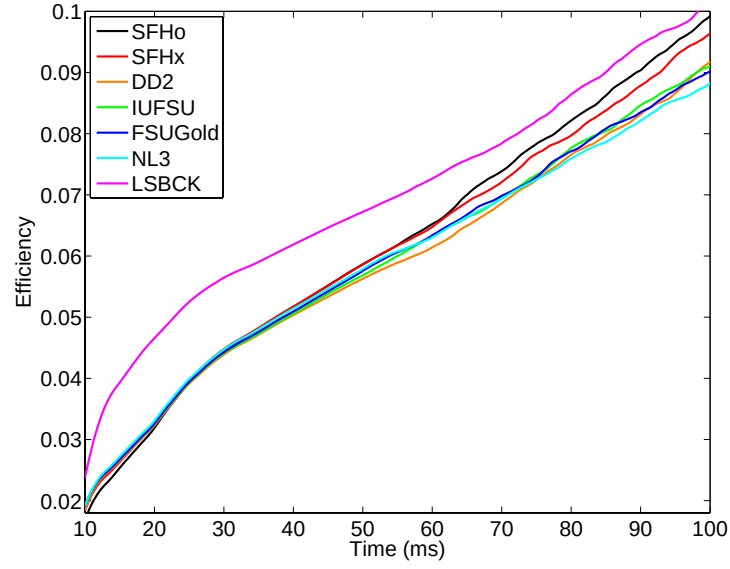


Figure 5.51: Heating Efficiency vs Time, 100 ms

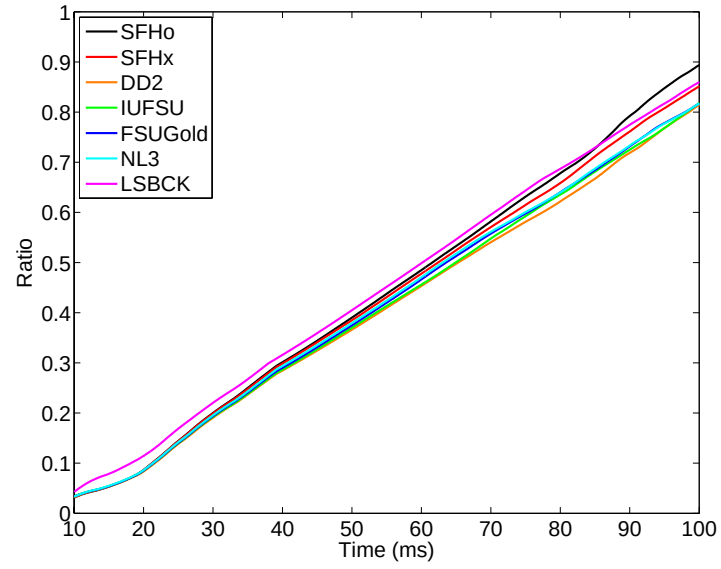


Figure 5.52: τ_{adv}/τ_{heat} vs Time

5.7 Effects of EOS by 200 ms

During the epoch between 100 and 200 ms after bounce, a clear picture of a revived and expanding shock as well as the onset of explosion is found in all models, demonstrated by the mean shock radius in Figure 5.53. Additionally, the beginning of a consistent stratification among the models begins to become visible in the history of the accretion through the shock. In previous studies, the variation in the neutrino luminosities and RMS energies was shown to scale with the mass accretion rate into the PNS, $\dot{M}_{\text{PNS}}(t)$, and the growing PNS mass, $M_{\text{PNS}}(t)$ (Janka 2012; O’Connor and Ott 2013).

$$L_{\text{acc}} \propto (M_{\text{PNS}})(\dot{M}_{\text{PNS}}) \quad (5.5)$$

In Figure 5.56, the PNS mass of the LSBCK model can be seen to initially (~ 20 ms post bounce) lag behind the other models but by ~ 120 ms it becomes the most massive, due to this increased accretion, and continues to grow faster until 200ms. Likewise, increased PNS compactness correlates with increased luminosity and harder neutrino spectra (Figure 5.66) As can be seen in Figure 5.57 and Figure 5.58, the LSBCK PNS core becomes more compact (smaller PNS radius, higher central density) in this period, as it becomes more massive. An increased mass accretion rate and a higher PNS mass correlate with more compact cores (Suwa et al. 2013). Additionally, the parameterized study of Scheck et al. (2008) showed that the increased contraction rate of the PNS causes increased pdV heating of the PNS. This in turn causes the gravitational potential energy to be released as neutrino radiation more rapidly, which further leads to accelerated heating of the gain region. The PNS of the LSBCK model is contracting more rapidly than all other models. The contraction rates of the SFHo and SFHx models are higher than the other RMF models, forming an intermediate group between LSBCK and the others (Figure 5.57). Likewise, the increased compactness of the LSBCK model as well as those of SFHo and SFHx can be seen quantitatively in Figure 5.59.

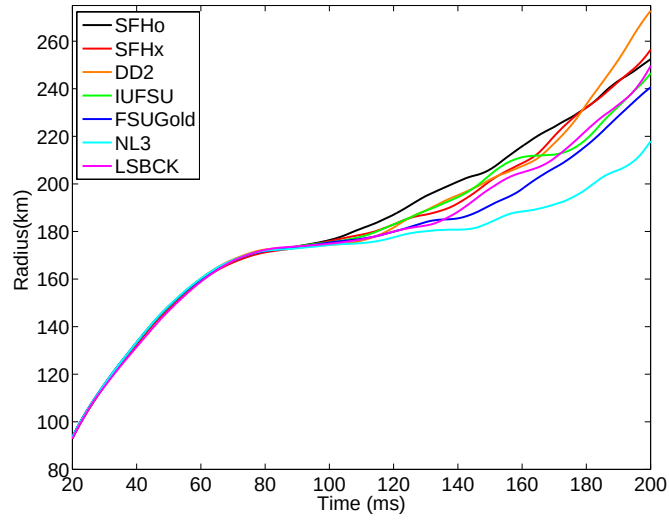


Figure 5.53: Shock Radius vs Time

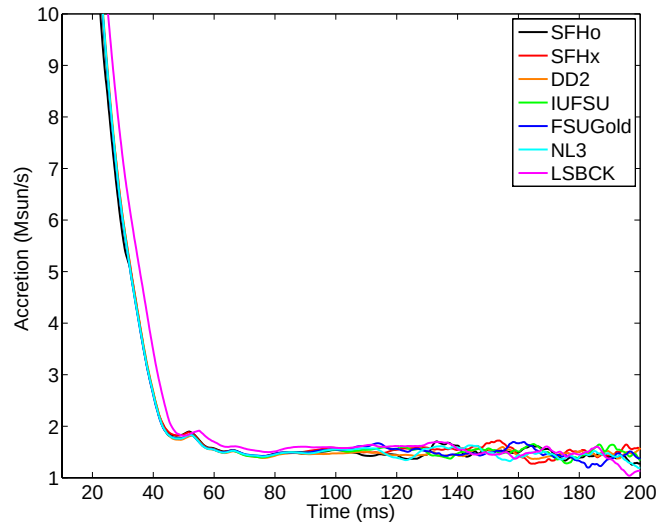


Figure 5.54: Gain Mass Accretion vs Time

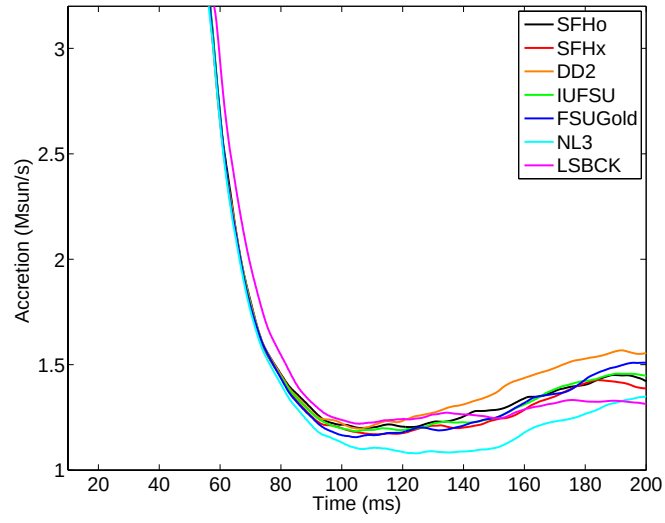


Figure 5.55: Shock Mass Accretion vs Time

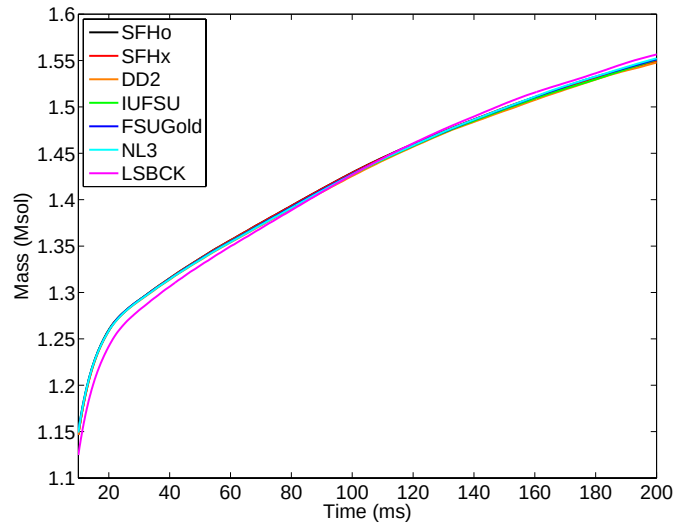


Figure 5.56: PNS Mass vs Time

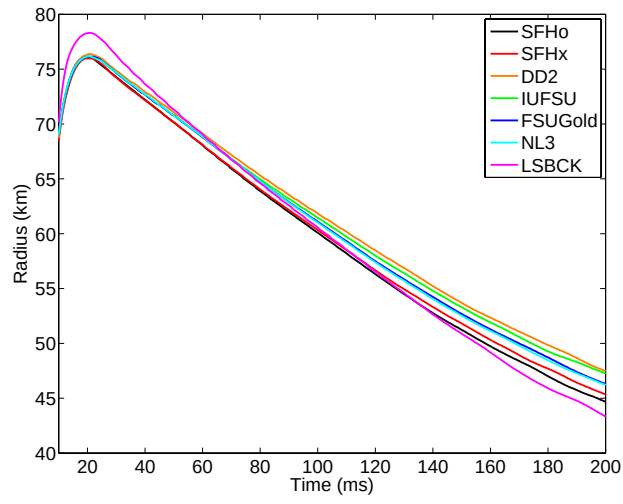


Figure 5.57: PNS Radius vs Time

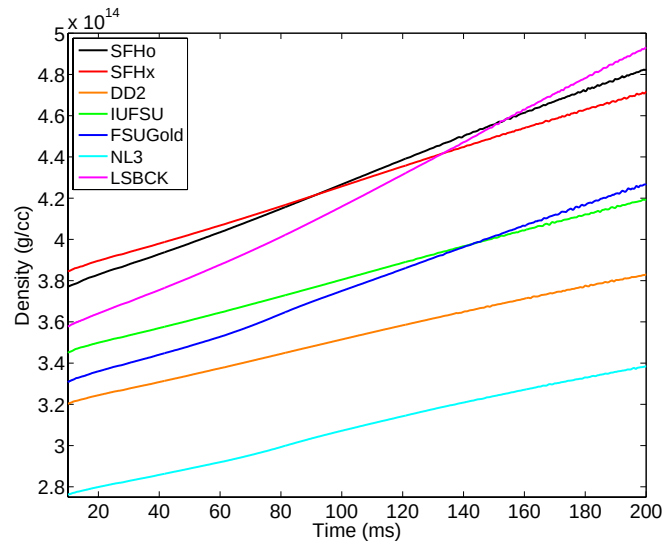


Figure 5.58: PNS Central Density vs Time

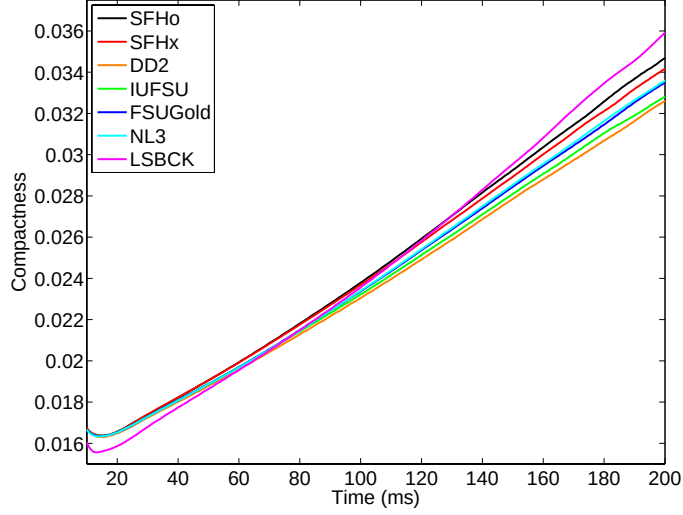


Figure 5.59: PNS Compactness vs Time. The LSBCK model’s PNS radius reaches the highest value among the models after rebounding from bounce, hence it has the lowest compactness at 20 ms. However as its radius contracts the most quickly among the models, its compactness overtakes that of the next most compact, the SFHo and SFHx models.

The luminosities of all neutrino species peak during this epoch, with the higher early accretion and higher mass of the LSBCK model correlating with its increased luminosities relative to the other models, with the NL3, SFHo, and SFHx models being the next most luminous (Figures 5.62, 5.63, 5.64, and 5.64). However, there is an interplay between the ram pressure of accretion and the boost to luminosity provided by the infalling/accreting material; higher accretion implies more inward ram pressure but also more fuel for mantle luminosity via accretion. Descending from the neutrino burst at shock breakout, the LSBCK, SFHx, and SFHo models all have relatively high ν_e luminosity. As accretion differences manifest in the ν_e luminosity peak after 60 ms, the LSBCK model becomes the most luminous for this neutrino species, as seen in Figure 5.62. The SFHx, and SFHo models have high $\bar{\nu}_e$ luminosity approaching the accretion peak. The LSBCK model becomes the most luminous for this neutrino species as well as for the $\nu_{\mu,\tau}$ and $\bar{\nu}_{\mu,\tau}$ (Figure 5.64, Figure 5.64), as its PNS becomes increasingly compact. The LSBCK model maintains this higher luminosity over the other models through 200 ms and beyond, as seen in Figure 5.63. The RMS energies of the SFHo and SFHx models are the highest in the early period, demonstrative of their hotter,

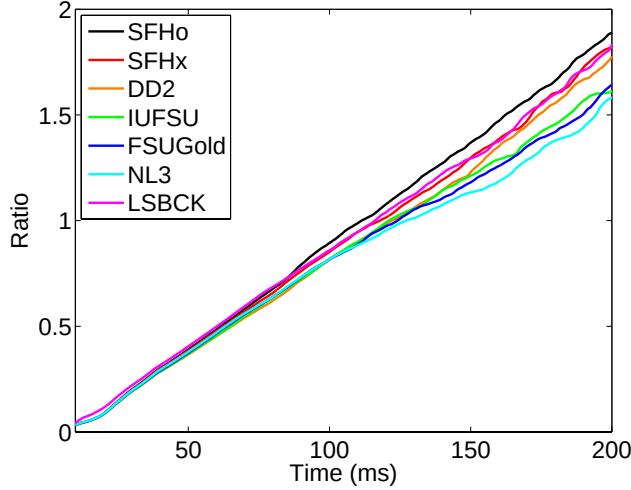


Figure 5.60: $\tau_{\text{adv}}/\tau_{\text{heat}}$ vs Time, 200 ms

more compact neutronospheres, though the RMS energy of the LSBCK model exceeds them once it exceeds them in compactness, as seen in Figure 5.66. Figure 5.67 shows the heating in the gain region by EOS model for bounce up to 200 ms. The LSBCK and SFHo EOS models show $\sim 20\%$ greater heating than the stiffer EOS group (IUFSU, NL3, FSUGold, and DD2) at 100 ms, extending to $\sim 25\%$ by 200 ms, with the SFHx EOS model falling in between. The timescale ratio and heating efficiency continue to rise monotonically, with SFHo, SFHx, and LSBCK higher than the other models, as seen in Figure 5.68 and Figure 5.69, respectively. Figure 5.71 and Figure 5.72 show that at 200 ms, the proton and neutron mass fraction of the LSBCK models is beginning to increase relative to the RMF models between 150 and 200 km, though it should be noted that at the relevant densities, the BCK EOS is the EOS being utilized.

By 200 ms after bounce, neutrino-driven convection has lead the shock front to have highly aspherical morphology in all models (Figure 5.74). This in turn has changed the character of the accretion, which is no longer generally radial but rather is beginning to be funneled by the aspherical shock into focused accretion streams. The focusing affect of the asphericity is due to the nature of how the shock interacts with the velocity of the infalling matter. As matter passes through the shock, the shock reduces the component of the infalling

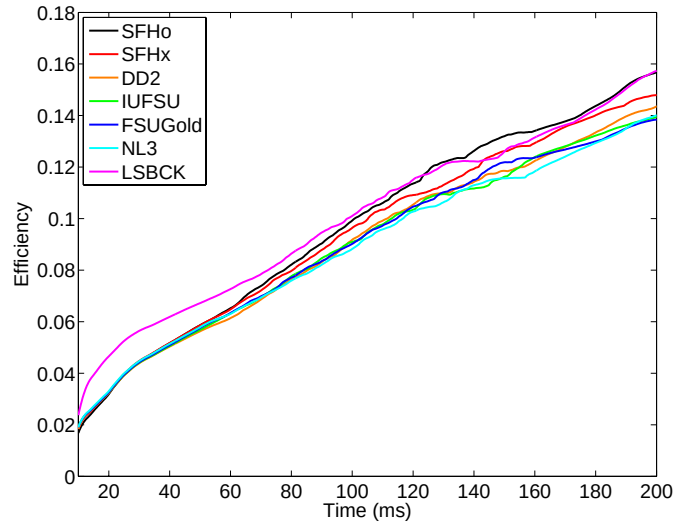


Figure 5.61: Heating Efficiency vs Time, 200 ms

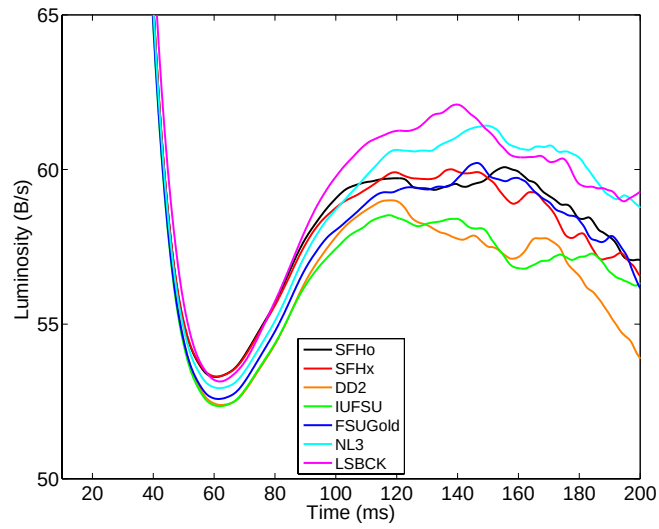


Figure 5.62: Electron Neutrino Luminosity vs Time, 200 ms.

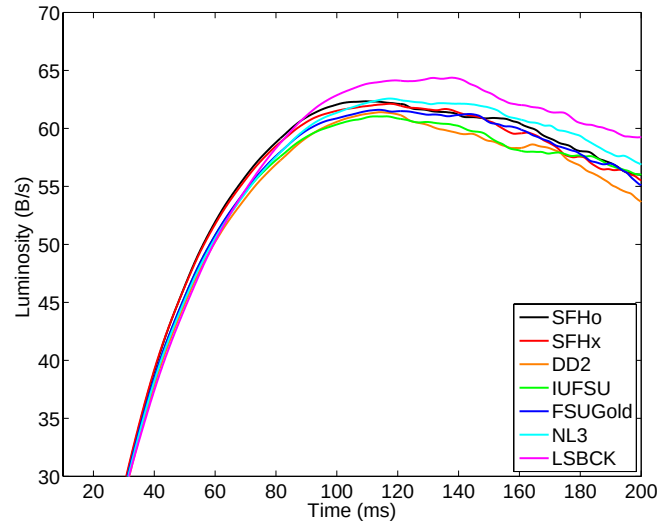


Figure 5.63: Electron Antineutrino Luminosity vs Time, 200 ms.

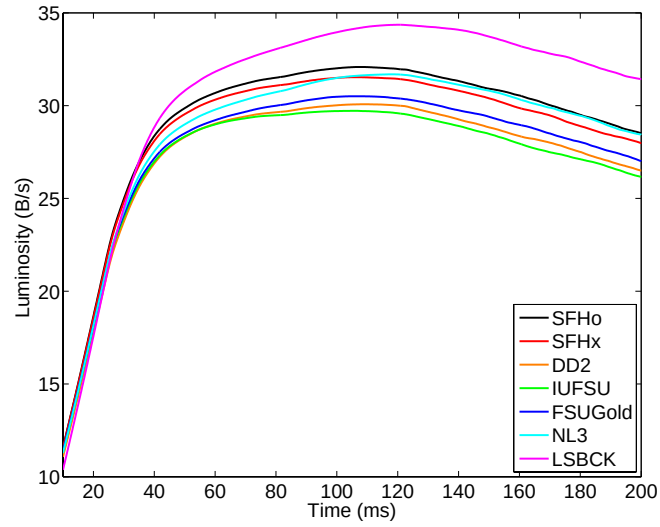


Figure 5.64: Mu and Tau Neutrino Luminosity vs Time, 200 ms

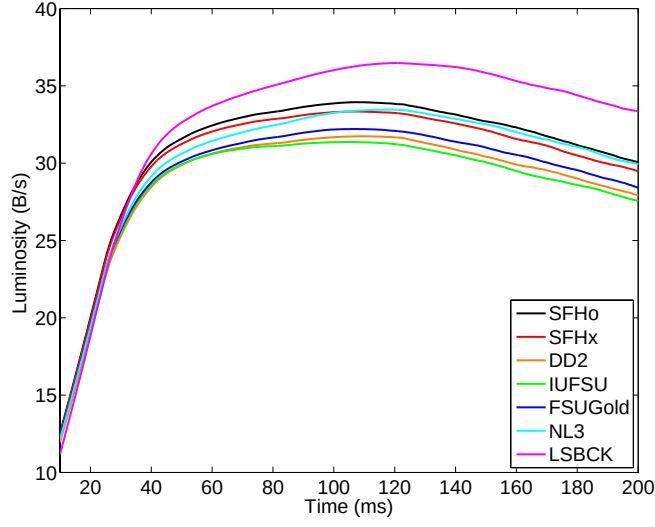


Figure 5.65: Mu and Tau Antineutrino Luminosity vs Time, 200 ms

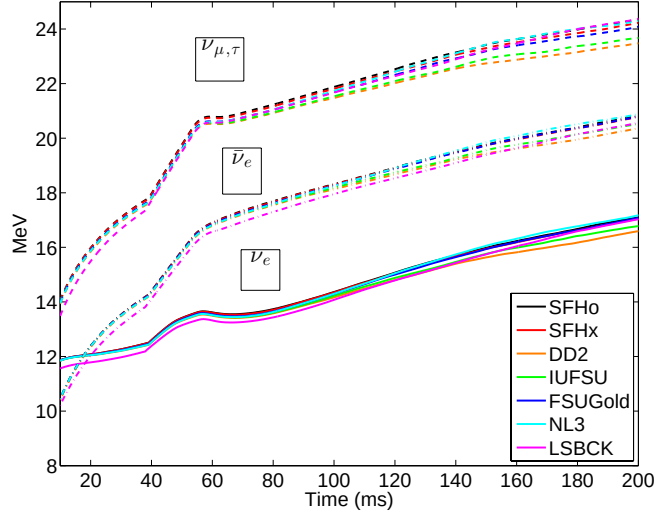


Figure 5.66: Neutrino RMS Energies vs Time, 200 ms.

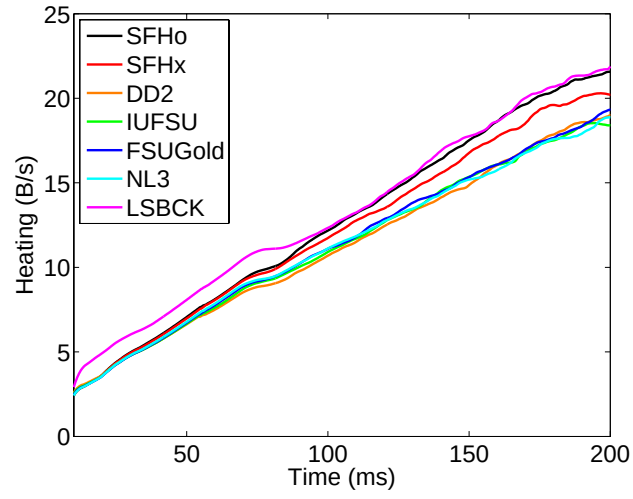


Figure 5.67: Gain Heating vs Time, 200 ms

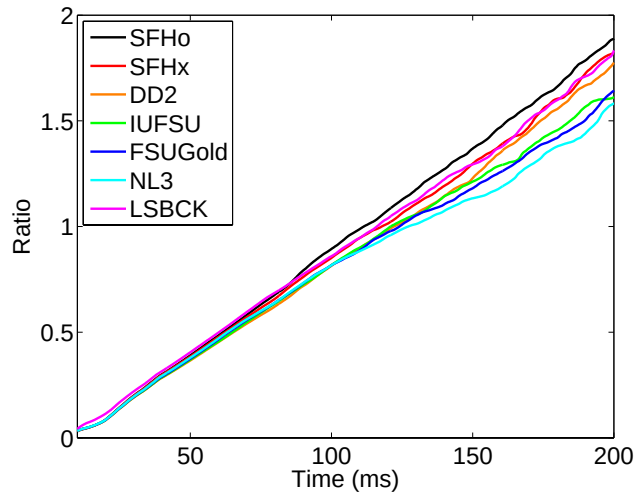


Figure 5.68: τ_{adv}/τ_{heat} vs Time, 200 ms

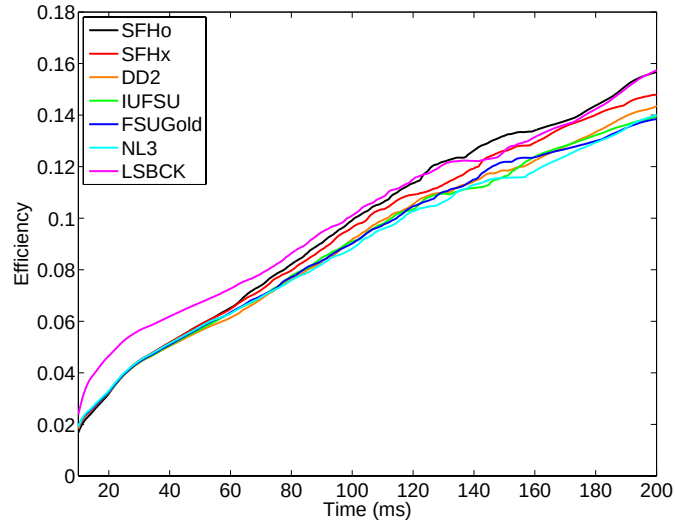


Figure 5.69: Heating Efficiency vs Time, 200 ms

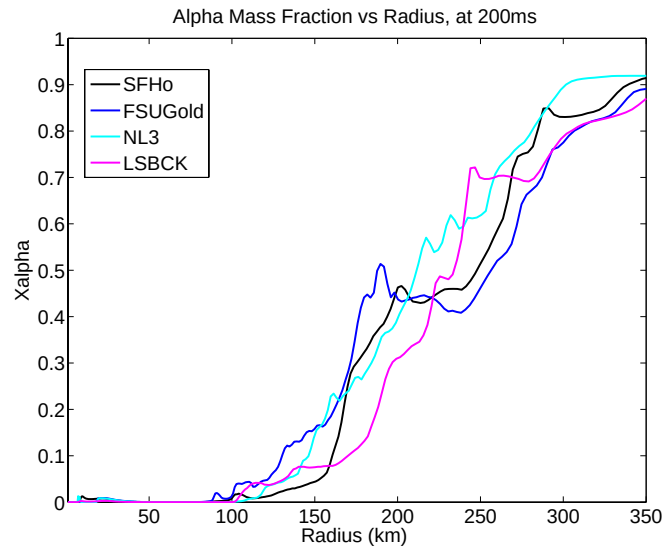


Figure 5.70: Alpha Mass Fraction vs Radius, 200 ms

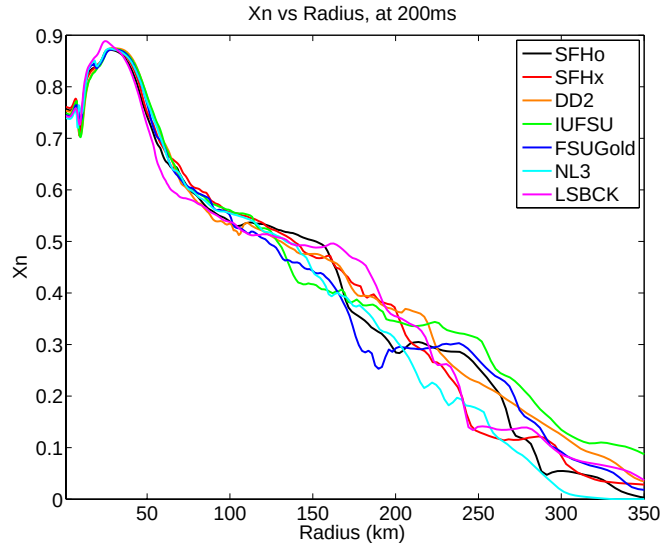


Figure 5.71: Neutron Mass Fraction vs Radius, 200 ms

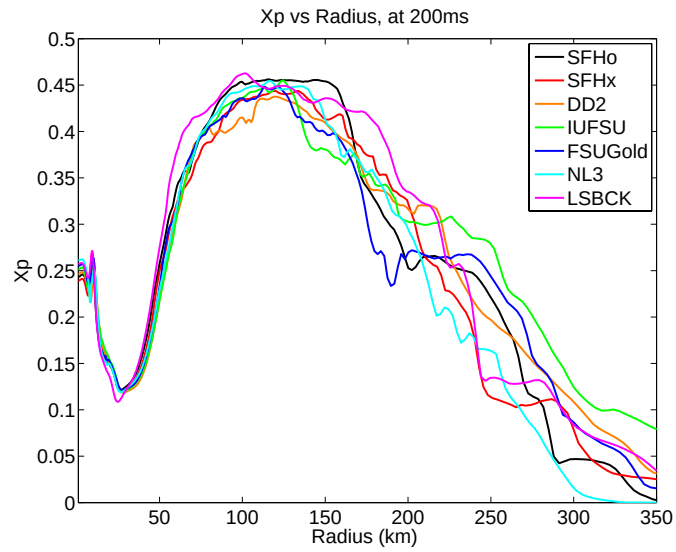


Figure 5.72: Proton Mass Fraction vs Radius, 200 ms

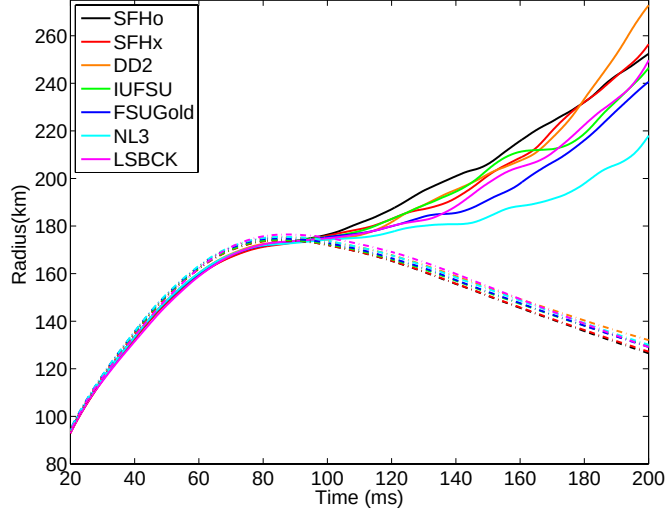


Figure 5.73: Shock Radius vs Time. The two-dimensional models all undergo shock revival just after 100 ms, while the shock radii of the one dimensional models monotonically contracts for the remainder of the simulation.

matter’s velocity that is normal to the shock, but not the component that tangential to the shock. This leads to a lesser reduction in the infall velocity of matter entering the shock where the shock surface is closer to tangential relative to the radial direction. As described in [Bruenn et al. \(2016\)](#), this leads to the formation of a post-shock boundary layer of such material which effectively serves to funnel accreting material from regions with large shock radii to regions where the shock is closer to the PNS. This can lead to the post-shock flow becoming supersonic. The result of this focusing of high-velocity accretion by aspherical shock morphology is that less kinetic energy of the infalling matter is thermalized, because material in such flows has lost less velocity to the jump conditions at the shock. This boosts the amount of additional kinetic energy these streams can deposit via impact on the PNS, where the kinetic energy is now thermalized, heating the PNS mantle and further boosting the neutrino luminosity. This accretion-stream-focusing effect can be seen in all of these models. Quasi-equatorial accretion streams occur for all models, with the a boundary layer of post-shock, low-entropy accreted material visible in each model (the material that appears aqua to green in Figure 5.74); material funneled from the expanded polar shock towards the compact equatorial shock region and then down to the PNS while still having relatively

low entropy due to this effect. Figure 5.75 shows a central equatorial accretion stream is approaching the PNS at 15% the speed of light, with a mach number $\mathcal{M} > 2$, in the SFHx model. Additionally, in Figure 5.69, we can see the effect of these accretion streams on the heating efficiency. The efficiencies continue upwards as convective activity in the gain region benefits from this enhanced delivery of this cooler material downward, where it can more effectively absorb neutrino energy.

5.8 Effects of EOS at 300ms

Between 200 and 300 ms after bounce, energy from neutrino heating continues to be deposited in the turbulent convection region below the shock and the shock accelerates rapidly, as seen in Figure 5.76, similar to Bruenn et al. (2013, 2016). Whereas the shock radius evolution for all models had up to 200 ms been roughly similar, in this accelerating phase we begin to see the radial shock trajectories diverge, with the “soft” group of equations of state, LSBCK, SFHo, and SFHx, accelerating more strongly. After 200 ms, the luminosities of all neutrino species decline steadily from the peak and are continuing this trend at 300 ms after bounce. In this epoch, the first appearance of non-zero diagnostic energy for all models can be seen (Figure 5.77), signifying that the onset of the supernova explosion has begun.

The $\tau_{\text{adv}}/\tau_{\text{heat}}$ ratio continues to rise for all models, as seen in Figure 5.78, most steeply for the SFHo, LSBCK, and SFHx models. The heating efficiencies η_{heat} (see Figure 5.79) begin to level off in this period at between 16 to 22%, with the soft group also being the most efficient. Also during this period, the gain volume begins to inflect upward, as seen in Figure 5.80, mirroring the trends in the shock radius and diagnostic energy, with SFHo, SFHx, and LSBCK being the leading group. Also, The LSBCK model has the highest mass neutron star for the majority of this period (Figure 5.81), the result of earlier increased mass accretion through the shock. it is in this epoch that by ~ 275 ms after bounce, the LSBCK PNS is no longer the most massive (Figure 5.81), with the NL3 and FSUGold models obtaining higher masses. Simultaneous with these effects (Figure 5.82) each model contains

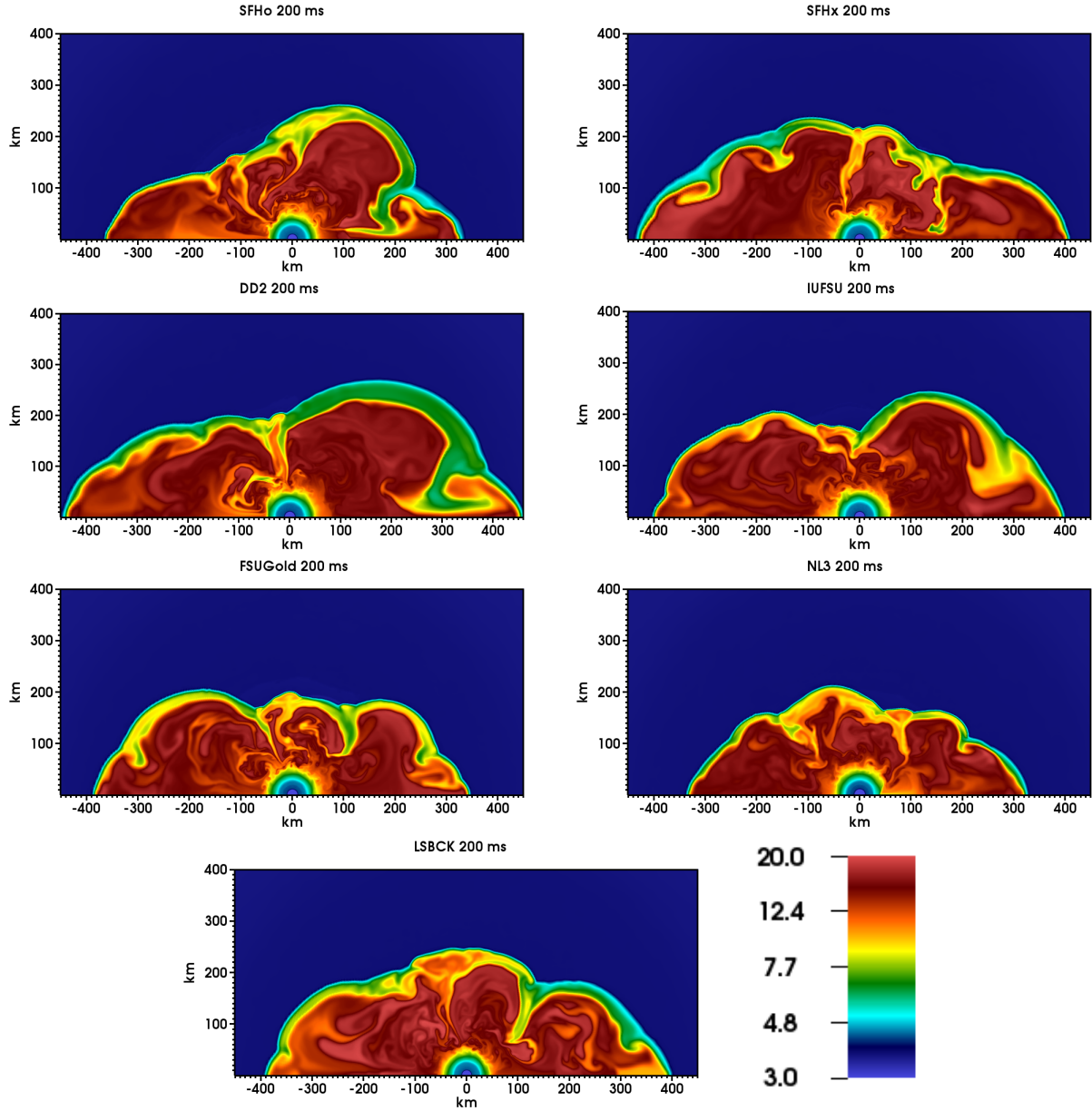


Figure 5.74: Entropy (log scale), 200 ms

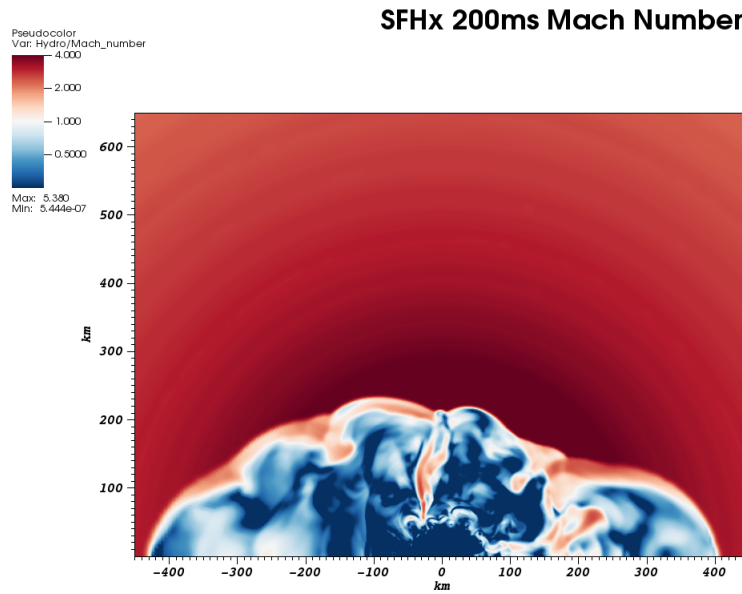
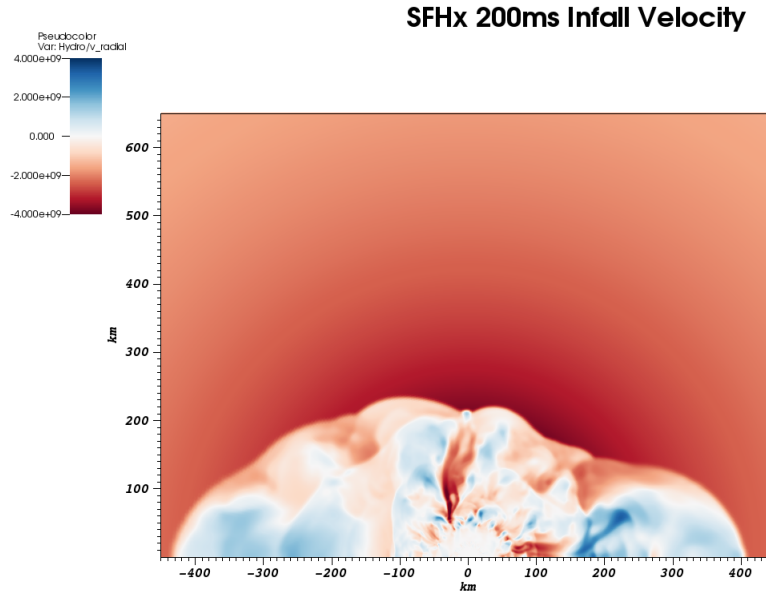


Figure 5.75: SFHx Accretion Velocity and Mach Number, 200 ms

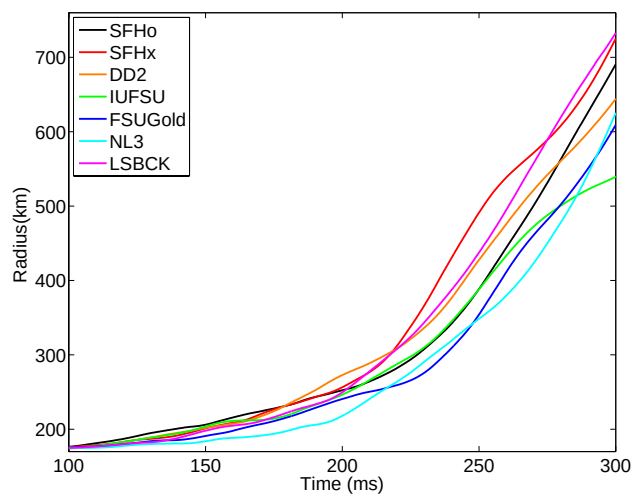


Figure 5.76: Shock Radius vs Time, 300 ms

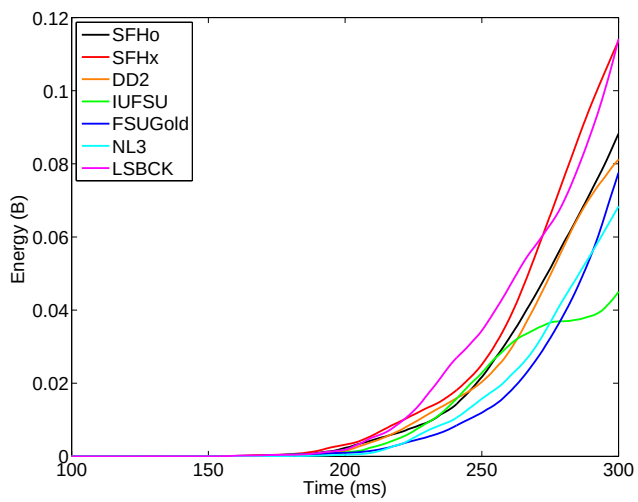


Figure 5.77: Diagnostic Energy vs Time, 300 ms

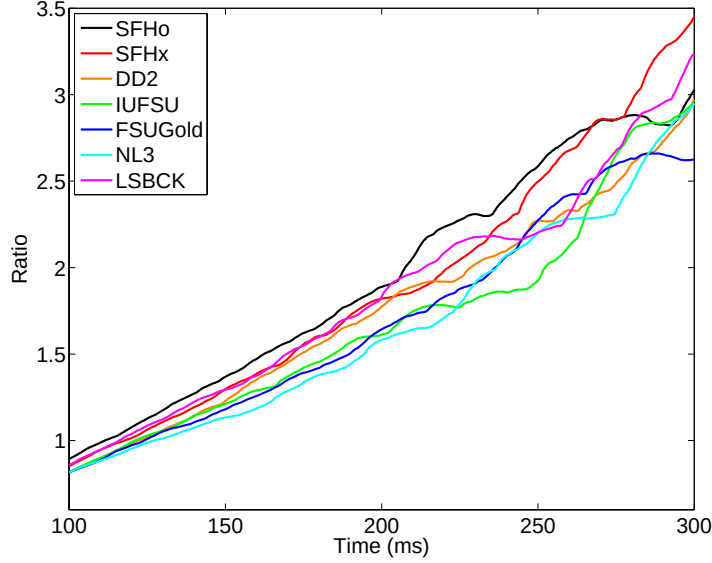


Figure 5.78: Timescale Ratio vs Time, 300 ms

one or more equatorial accretion streams in the compact shock region and a 2:1 elongation towards the pole. The relatively lower mass fraction of alpha particles in the LSBCK model relative to the others (Figure 5.83) leads to an increased neutrino loss in the non-LSBCK models, while the proton mass fraction excess for LSBCK leads to more available targets for heating by $\bar{\nu}_e$ captures on protons (Figure 5.84).

5.9 Effects of EOS at 500ms

At 500 ms after bounce, the diagnostic energy of each model is rising monotonically, with the LSBCK and SFHo models rising the most steeply, with the LSBCK model roughly twice as energetic (Figure 5.3). The shock radius evolution for all models continues to diverge (Figure 5.2), with LSBCK, SFHo, and SFHx increasing at a higher rate. The ν_e and $\bar{\nu}_e$ luminosities all continue to decline (Figure 5.5 and Figure 5.6) as well as those for the ν_μ and ν_τ and their antineutrinos (Figure 5.7 and Figure 5.8) in tandem with the gain mass accretion rate, as shown in Figure 5.10. The ν_e and $\bar{\nu}_e$ luminosities are similar to $\dot{M}_{PNS}$ in this phase but not so for the ν_μ and ν_τ and their antineutrinos.

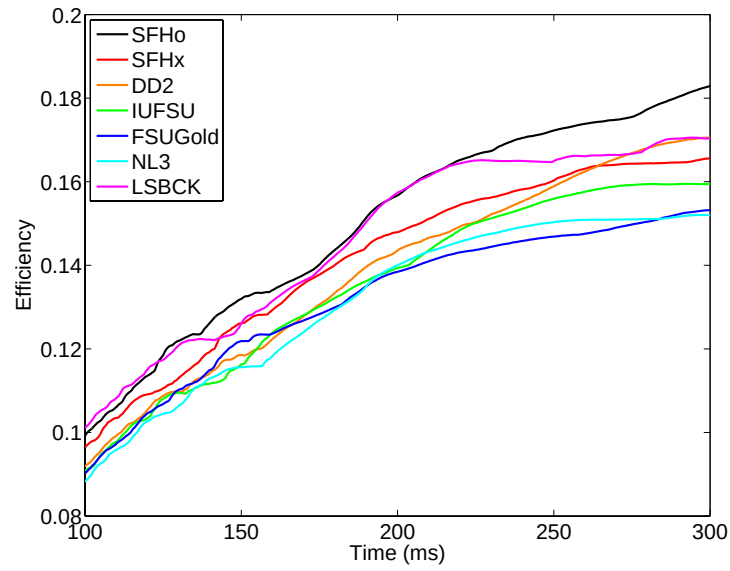


Figure 5.79: Heating Efficiency vs Time, 300 ms

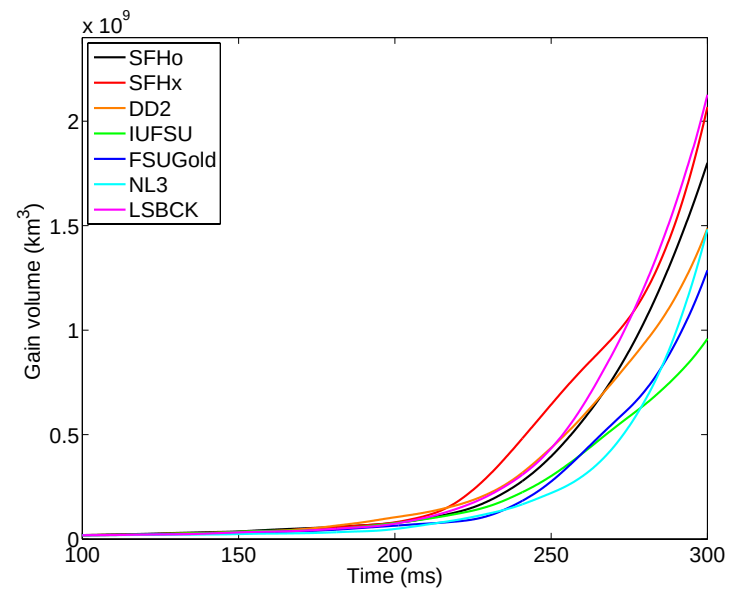


Figure 5.80: Gain Volume vs Time, 300 ms

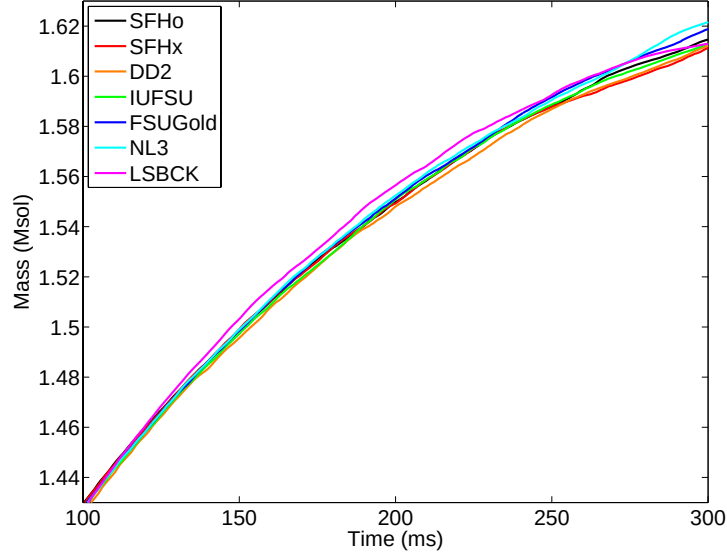


Figure 5.81: PNS Mass vs Time, 300 ms

The shock mass accretion rate becomes negative for all models between 300 and 400 ms, with the rates for the trio of LSBCK, SFHo, and SFHx going negative at ~ 340 ms (see Figure 5.85). The shock is now beginning to lift the envelope away from the star at roughly $1M_{\odot} \text{ sec}^{-1}$, a value that will gradually fall as the envelope density and shock velocity decline. Note that while the net shock accretion has become negative, the accretion to the gain radius still remains positive, as shown in Figure 5.10. From this point, though accretion continues onto the PNS, accretion comes from matter already in the hot bubble.

Several features in the luminosity become apparent in this period. An inflection in the rate of decline of luminosity for all neutrino species can be seen between 300 ms and 500 ms, gradually at first and more sharply after 400 ms. After 300 ms the steepness of all luminosities' decline lessens, and then just after 400 ms, the rate of decline of the ν_{μ} and ν_{τ} neutrino luminosities diminishes, with the ν_e and $\bar{\nu}_e$ luminosities for each model following suit ~ 50 ms later. At this point, the luminosities of all species for LSBCK not only lessen their decline but actually increase briefly, as do those of the ν_e and $\bar{\nu}_e$ for the NL3 model. Beyond this inflection, NL3 begins to have the sustained largest electron anti-neutrino luminosity of the group, followed by LSBCK, while LSBCK has higher ν_{μ} and ν_{τ} luminosities. The

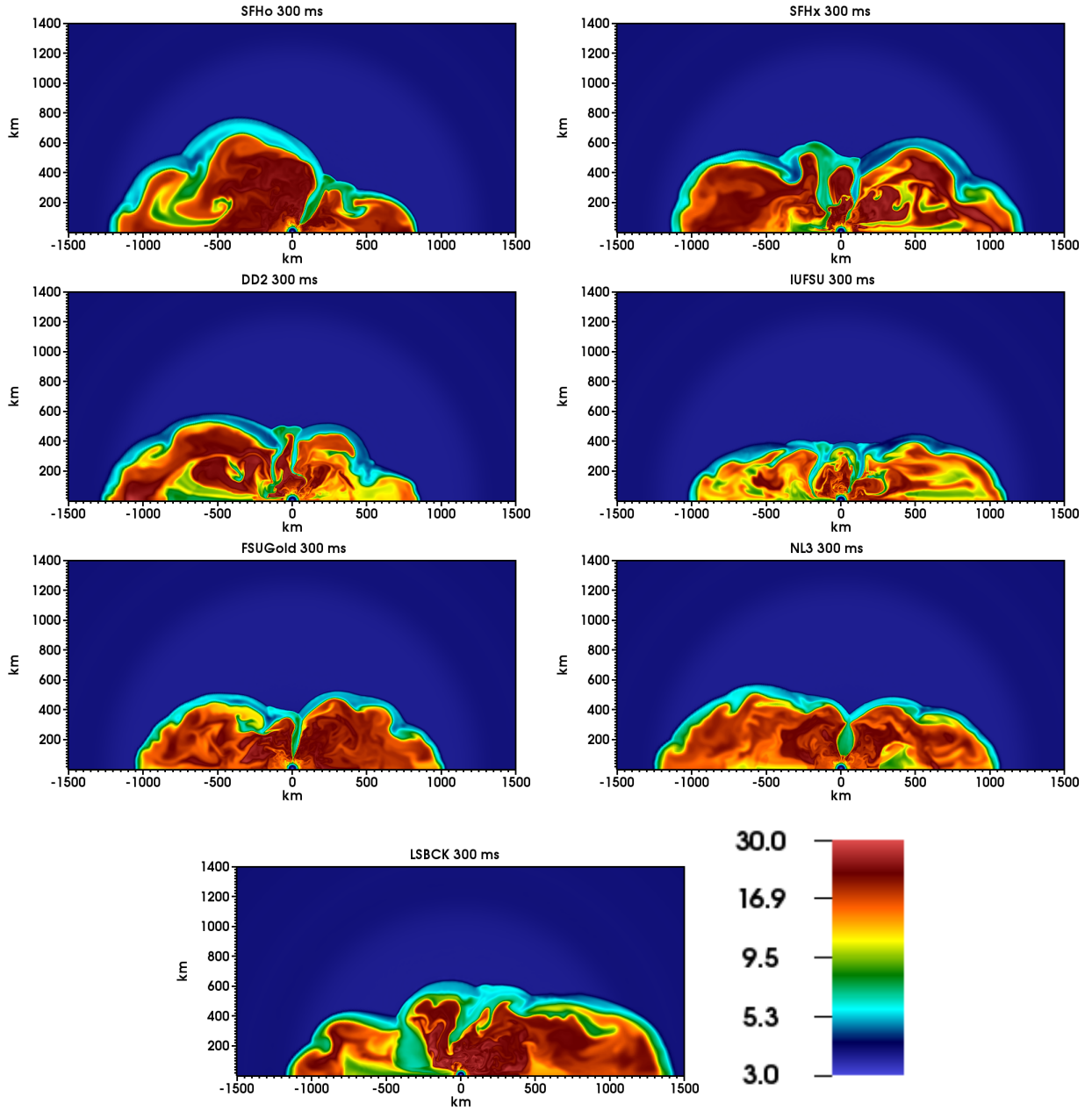


Figure 5.82: Entropy (log scale), 300 ms

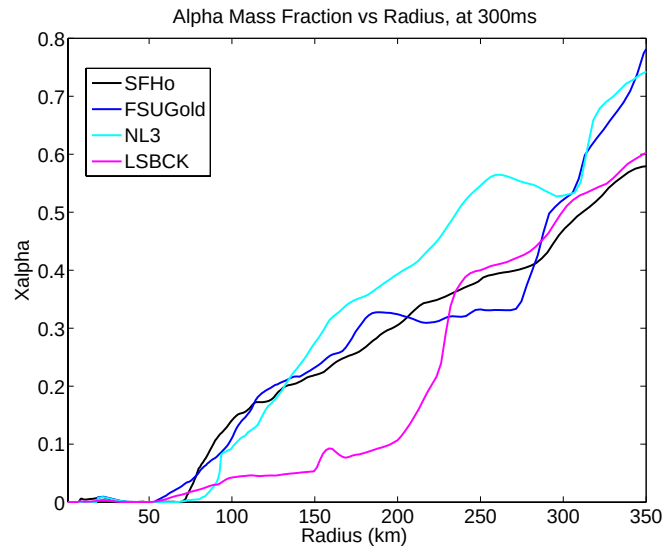


Figure 5.83: Alpha Mass Fraction vs Radius, 300 ms

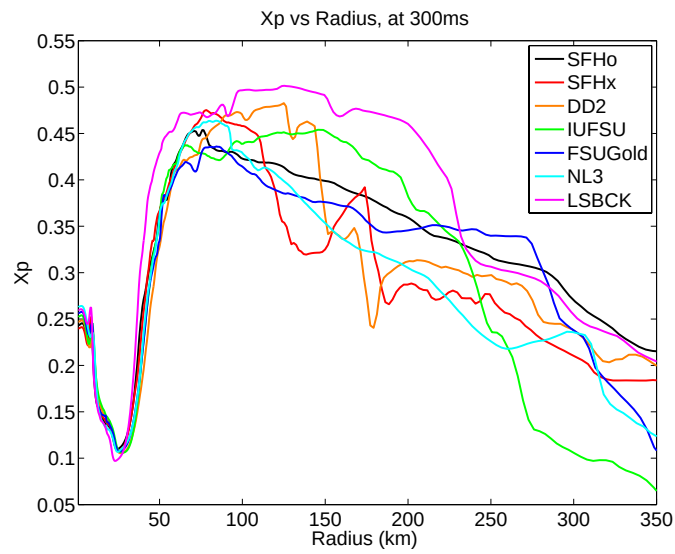


Figure 5.84: Proton Mass Fraction vs Radius, 300 ms

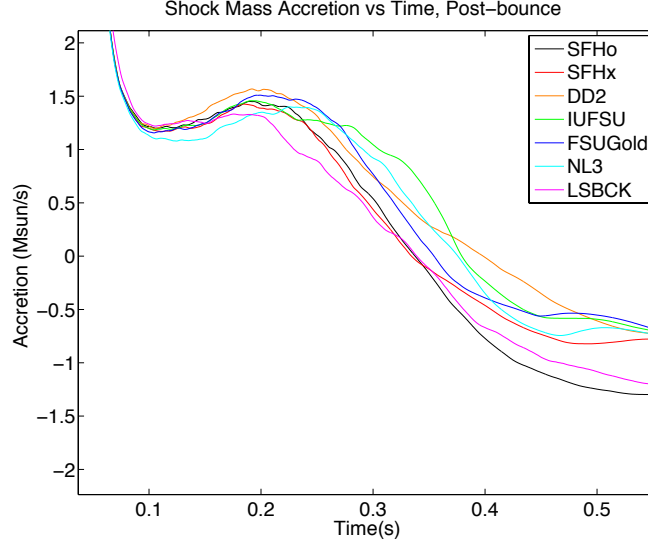


Figure 5.85: Shock Mass Accretion vs Time, 500 ms

luminosity up to this point has been dominated by an accretion-driven “mantle” luminosity component (as discussed in section 5.7), with a smaller “core component” (Suwa et al. 2016), characterized by the neutrino emission from the PNS. As the net accretion has declined, the change in the shape of the luminosity terms reflects the transition to a period where the core component (characterized by the PNS formation history during collapse and bounce, prior to accretion) now more strongly influences the luminosity profile. This transition is referred to in Suwa et al. (2016) as the “turning point”, and they discussed this transition as a diagnostic of explosion onset, though the models discussed in this dissertation have all been exploding for 200 ms or more at this point, indicating that changes in accretion should be a trailing indicator of explosion. Also, between 300 and 500 ms the PNS masses diverge further (Figure 5.86), and the “stiffer” DD2, IUFSU, FSUGold, and NL3 EOS models continue accrete mass to a greater degree, and hence their PNS masses continue to grow at a steeper rate.

As shown in Figure 5.87, the LSBCK model achieves the highest neutrino heating rate of over 20 B s^{-1} , followed by SFHo. Heating efficiency (Figure 5.88) peaks and then begins to decrease in this period, due to the decreasing density in the expanded gain region. The relatively lower mass fraction of alpha particles in the LSBCK model relative to the

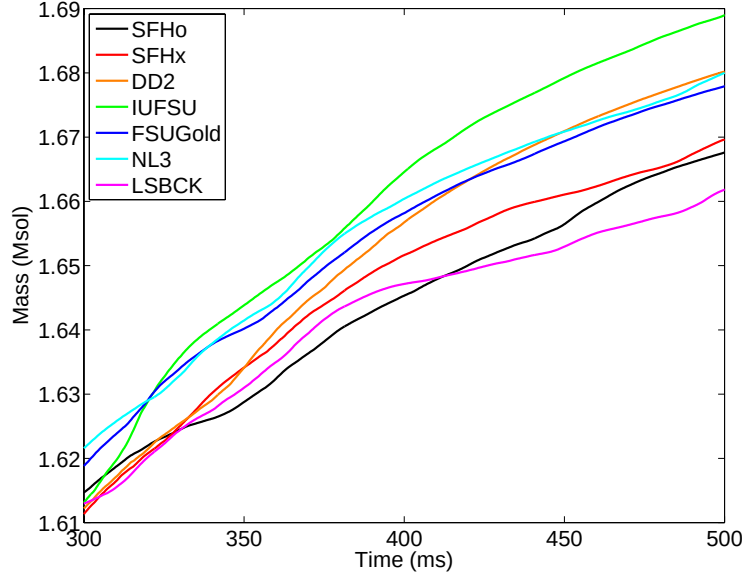


Figure 5.86: PNS Mass vs Time, 500ms

others (Figure 5.89) leads to an increased neutrino loss in the non-LSBCK models. We also see a proton and neutron mass fraction excess for LSBCK, leading to more available targets for ν_e captures on neutrons (Figure 5.90) and $\bar{\nu}_e$ captures on protons (Figure 5.91). Simultaneous with these effects (Figure 5.92) shock expansion continues, with sustained equatorial accretion streams in the compact shock region and a sustained 2:1 elongation towards the pole.

5.10 Effects of EOS at 1 second

At 1 second after bounce, the shock front radius of all models is expanding nearly linearly as shown in Figure 5.93, with the LSBCK model maintaining the steepest rate of shock expansion, with a mean shock radius of ~ 9400 km. Figure 5.3, shows the leveling of the diagnostic energy for each model, and also shows that the models are at different points in the development of their explosion. SFHo is in the most complete explosion phase, with the diagnostic energy leveling off, followed by SFHx, FSUGold, IUFSU, and DD2, with LSBCK and NL3 still in a less complete phase, with their diagnostic energy still climbing as the

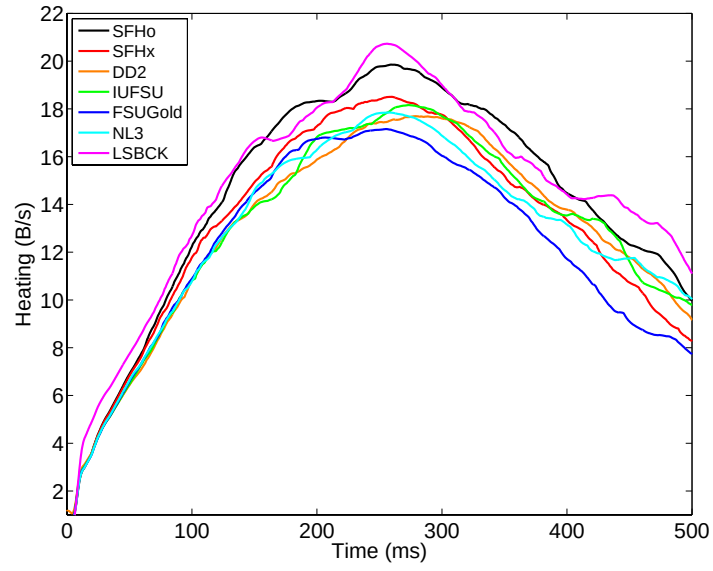


Figure 5.87: Gain Neutrino Heating vs Time, 500ms

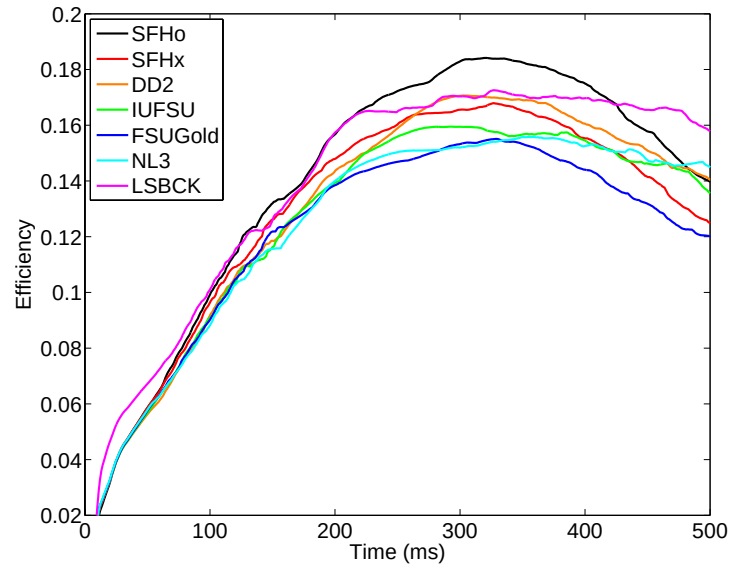


Figure 5.88: Heating Efficiency vs Time, 500ms

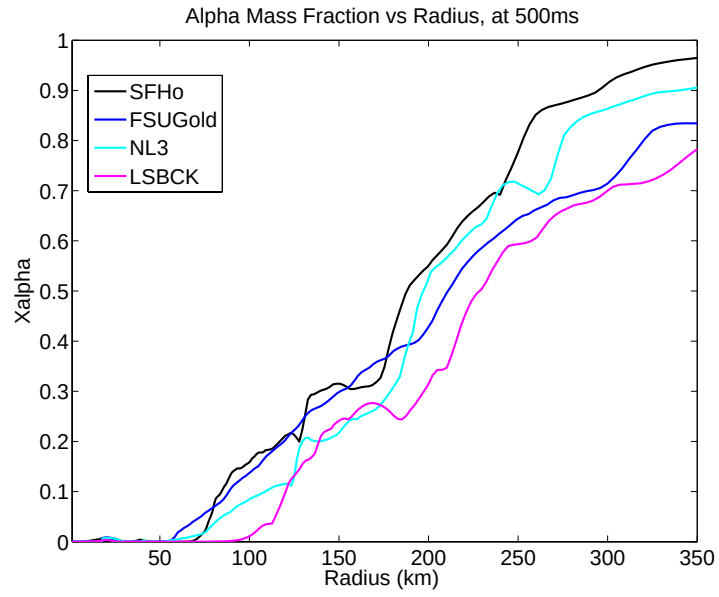


Figure 5.89: Alpha Mass Fraction vs Radius, 500ms

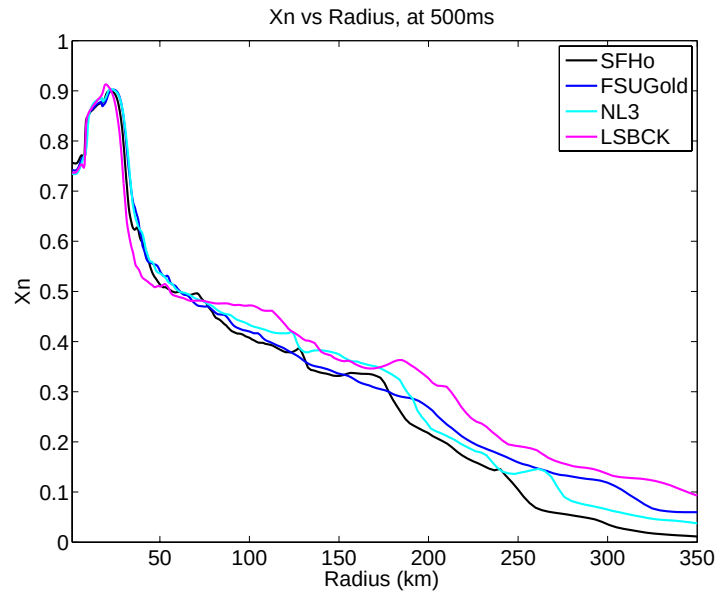


Figure 5.90: Neutron Mass Fraction vs Radius, 500ms

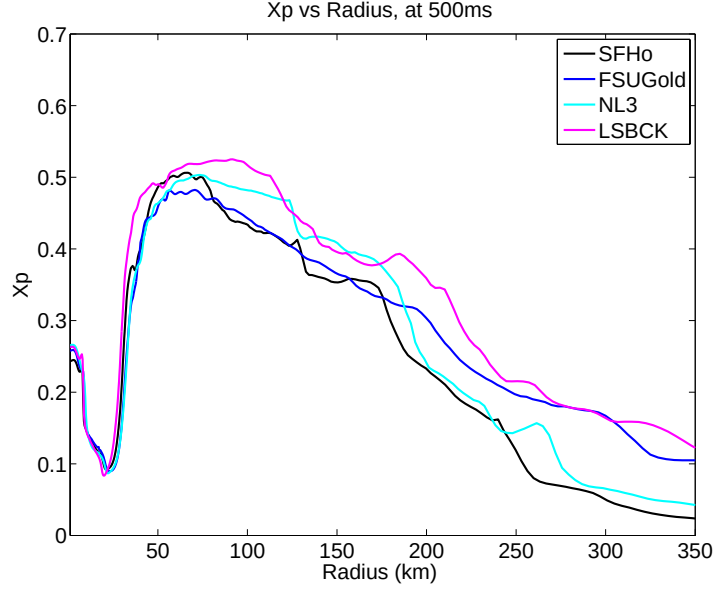


Figure 5.91: Proton Mass Fraction vs Radius, 500ms

energetic motivation of their explosions is still taking place. The FSUGold model is the least energetic explosion, in tandem with having the smallest shock radius (~ 7200 km) at 1 second after bounce. At this point in the simulations, SFHo has the second highest shock radius at just over 8000 km, followed closely by the NL3 and IUFSU models. In this it can be seen that the ranking of explosion energies does not correlate exactly with the current shock radii, in that several of the models are still evolving.

The LSBCK and NL3 models, as in Figure 5.94, have explosion geometries with a single accretion stream. The LSBCK model has a single polar accretion stream spanning the whole southern hemisphere and a massive heated upflow on the other pole. The NL3 model explodes via two polar heated plumes fed by a wide equatorial accretion stream.

5.11 Neutrino Heating by Individual Processes

The post-processing calculation of the neutrino heating the heating of a gain layer zone by multiplying the zone density by the zone volume and then by the neutrino energy deposition

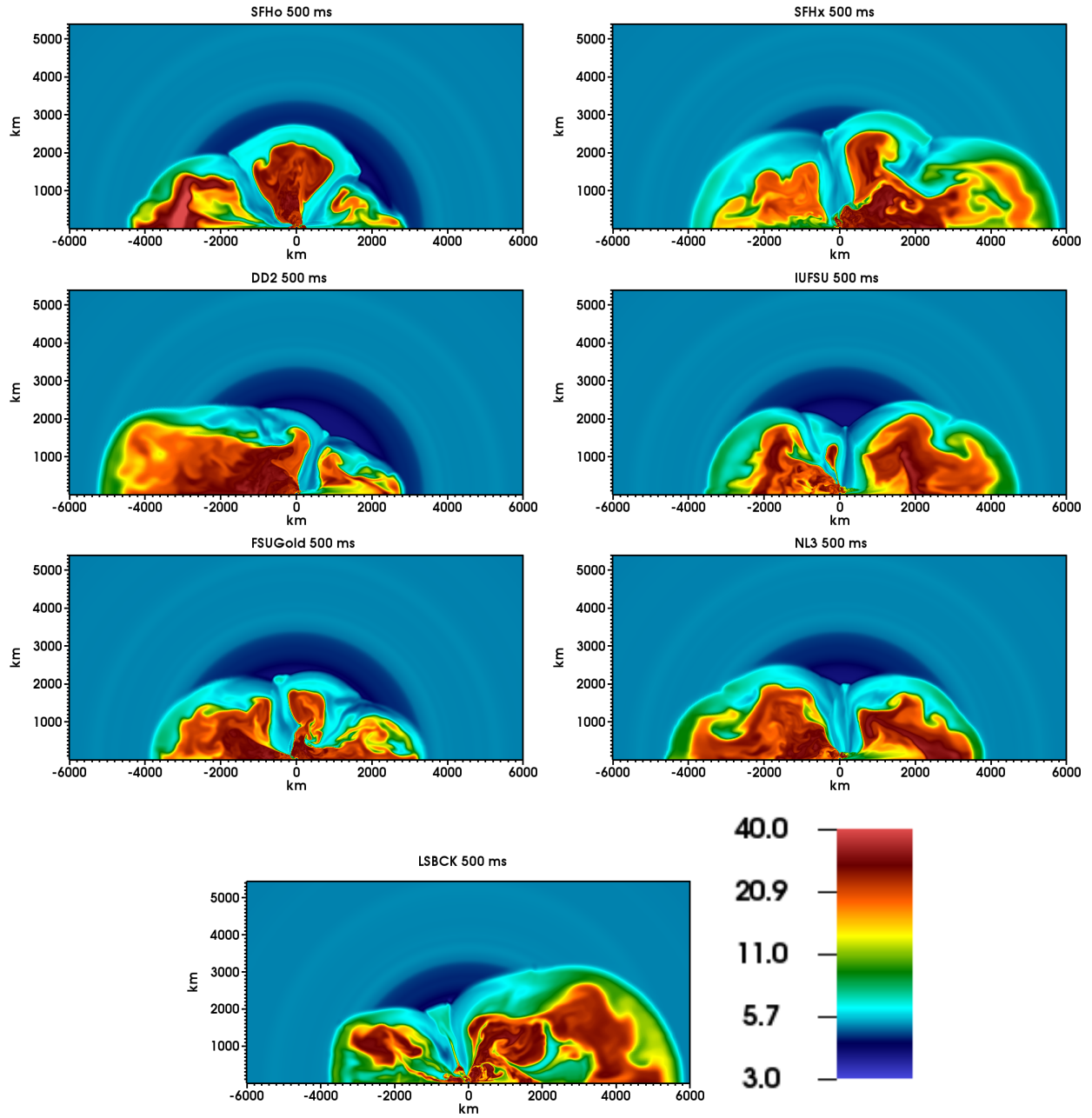


Figure 5.92: Entropy (log scale), 500 ms

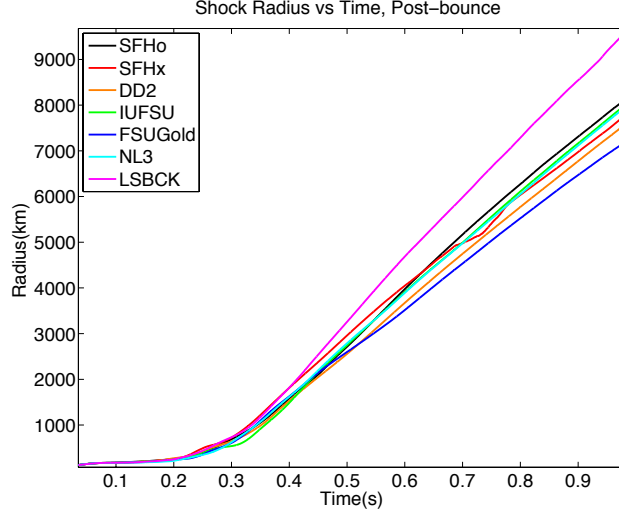


Figure 5.93: Average Shock Radius vs Time, 1 second after bounce. As many of these models are highly aspherical at late times, the average shock can be significantly less than the maximum value of the shock radius.

rate, but does not isolate the effects of composition nor for heating process, i.e.

$$\nu_e + n \rightleftharpoons p^+ + e^- \quad (5.6)$$

$$\nu_{\bar{e}} + p \rightleftharpoons n + e^+ \quad (5.7)$$

$$\nu_e + e^- \rightarrow \nu_e + e^- \quad (5.8)$$

$$\nu_{\bar{e}} + e^- \rightarrow \nu_{\bar{e}} + e^- \quad (5.9)$$

$$\nu_e + n \rightarrow \nu_e + n \quad (5.10)$$

$$\nu_e + p \rightarrow \nu_e + p \quad (5.11)$$

and the μ and τ equivalents, which depend on composition. Composition and heating by interaction process is output by CHIMERA, but on a radial ray basis, only for every few thousand timesteps, and is not realistically extractable via post-processing from multidimensional simulation runs at the moment. However, investigating such outputs from the only radial ray of a 1D simulation, before the behavior of 1D and 2D simulations diverge,

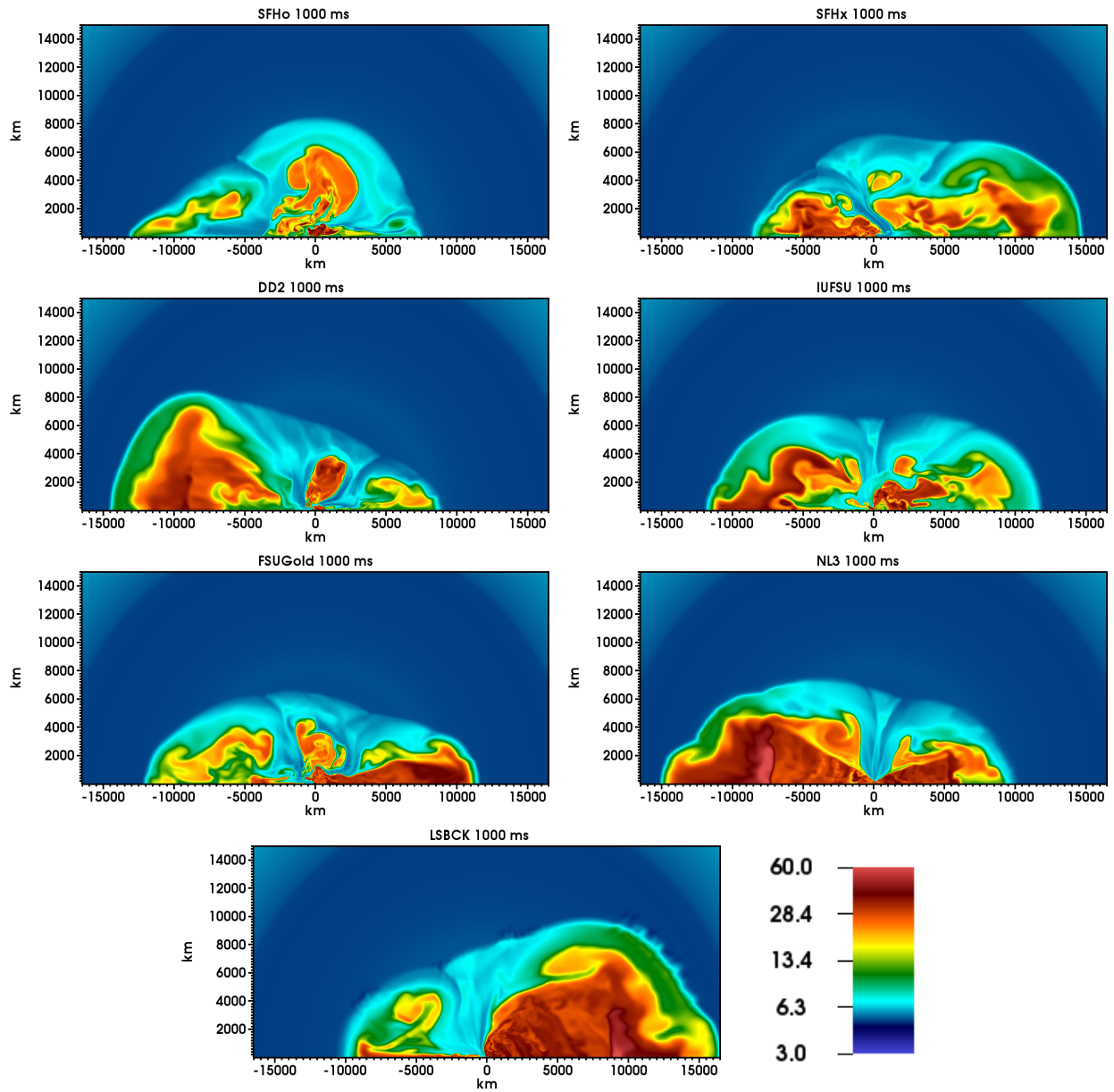


Figure 5.94: Entropy (log scale), 1 second

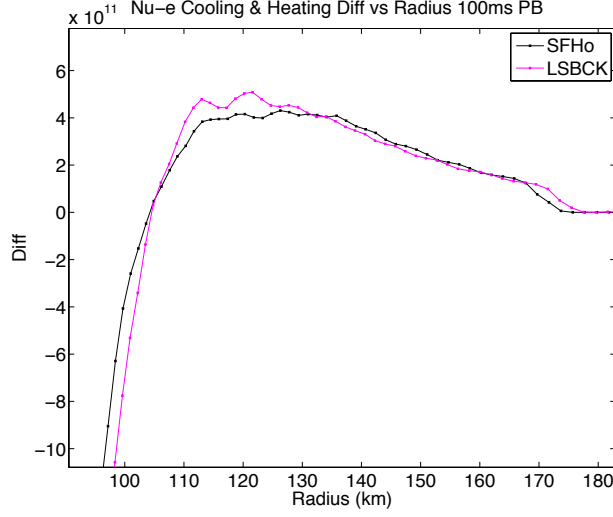


Figure 5.95: Electron Neutrino Net Heating by Absorption/Emission vs Radius, 100 ms

demonstrated that the composition and net heating by absorption/emission made up to a $\sim 25\%$ difference in favor of LSBCK for ν_e (Figure 5.95) and a $\sim 25\%$ difference for $\bar{\nu}_e$ (Figure 5.96). When composition at 100 ms is taken into account, the SFHo model has an excess of protons in the gain layer compared to the LSBCK model (Figure 5.97), with the reverse being true with regards to neutrons (Figure 5.98). This means that there are more potential targets in the SFHo model for heating by $\bar{\nu}_e$ captures on protons, and the LSBCK model being favored to a lesser extent for ν_e captures on neutrons, in the case of equal luminosity between the models. Even with SFHo being favored for heating by $\bar{\nu}_e$ captures, Figure 5.95 shows that the luminosity advantage of the LSBCK model overcomes this and leads to more net heating in the LSBCK model. Like in the early accretion epoch, the increased alpha fraction leads to a reduction in net neutrino opacity. This implies that where we see a near equivalence in heating rates of SFHo and LSBCK for the epoch of 100 ms to 200 ms in Figure 5.67, the heating rate of LSBCK is most likely underestimated relative to SFHo by an appreciable fraction.

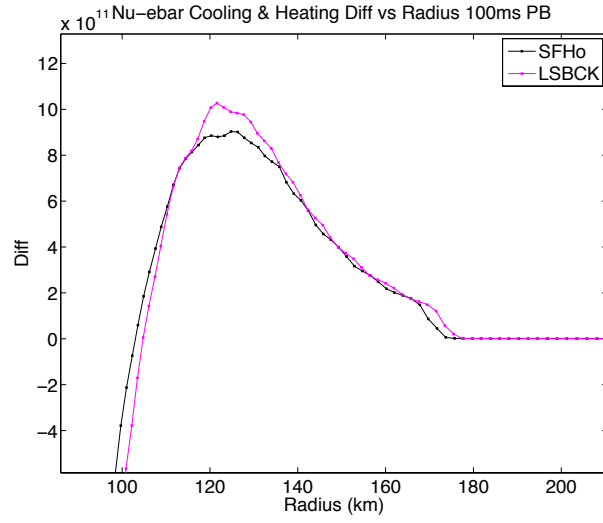


Figure 5.96: Electron Antineutrino Net Heating by Absorption/Emission vs Radius, 100 ms

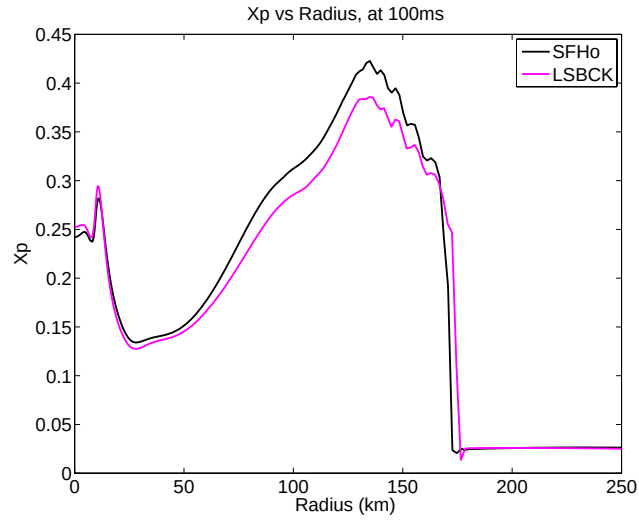


Figure 5.97: Proton Mass Fraction vs Radius, 100 ms

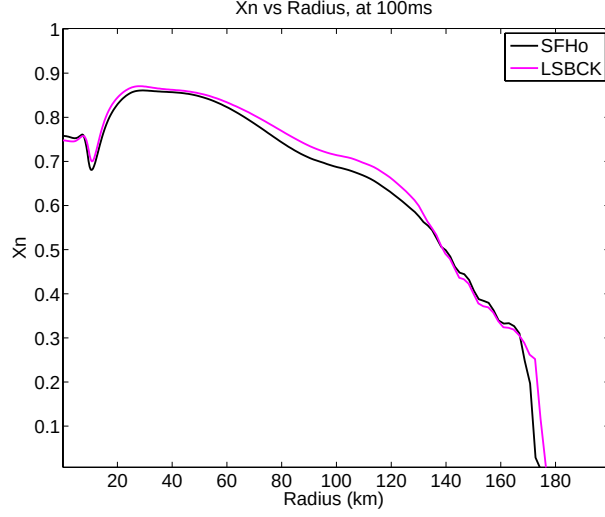


Figure 5.98: Neutron Mass Fraction vs Radius, 100 ms

5.12 Summary

5.12.1 EOS Effects on Neutrino Luminosities & Energies

The effect of the EOS on the luminosity profile of each model begins at the outset of each simulation, with the character of collapse being influenced by the different composition and thermodynamic state the collapsing core material. The composition affects the neutrino luminosity behavior during collapse, as discussed in Section 5.2. Through the EOS, the deleptonization history of the core will influence both the formation and structure of the PNS as well as the evolution of the luminosity of each neutrino species over the course of the simulation.

As discussed in Section 5.3, the EOS leads to the different peaks of the electron neutrino luminosity between the different equations of state in the neutrino burst at shock breakout. The neutrino luminosities over time show the models maintaining a similar luminosity profile during the first 100 ms, followed by a hierarchy among the models following the luminosity peak of $2\text{--}5 \text{ B s}^{-1}$ (which occurs between 100–200 ms) of $2\text{--}5 \text{ B s}^{-1}$, and then roughly maintains this order until an additional divergence event occurs near 500 ms, as the heating

behavior transitions from being driven by the higher-accretion, mantle-luminosity plus core luminosity early epoch to the lower luminosity, more core-luminosity driven phase, where the luminosity profiles are more strongly characterized by the individual core properties, i.e. higher relative $\bar{\nu}_e$, $\nu_{\mu,\tau}$ and $\nu_{\mu,\tau}^-$, as can be seen in Figures 5.5 through 5.8. After 500 ms, the divergence of luminosities increases from 2–5 B s⁻¹ to up to 5–12 B s⁻¹, depending on the species. The luminosities of the $\bar{\nu}_e$ at peak during the convective period are slightly higher than for the ν_e , as is the case for the $\nu_{\mu,\tau}$ and $\nu_{\mu,\tau}^-$ as well. The downward mass accretion at the gain surface (Figure 5.10) reflects the evolution of the luminosity profiles in that while the LSBCK model has the highest gain accretion prior to 200ms, the resulting mantle luminosity translates to decreased gain accretion relative to the other models until the end of the mantle luminosity dominated epoch ~ 450 ms, beyond this point, the gain accretion of the LSBCK model begins to decline more slowly, becoming one of the models with the shallowest gain accretion decline (along with NL3, and for a time, DD2 and IUFSU), unlike the SFHo model, the gain accretion of which is the least of the models considered here. This translates into a shallower decline in the luminosities of all species for LSBCK and NL3, whereas the SFHo late neutrino ν_e and $\bar{\nu}_e$ luminosities and RMS energies are the least of the group. It should be noted, however, that the $\nu_{\mu,\tau}$ and $\nu_{\mu,\tau}^-$ luminosities of the SFHo model are relatively high, demonstrating that the emission profile of the heavier neutrino species is related to the behavior of the PNS (of which the SFHo model had among the steepest radial contraction rates, correlating with higher neutrino trapping), not the accretion-driven luminosity.

Initially in Figure 5.47 and in Figure 5.99, the effect of the EOS on the neutrino RMS energies can be seen. The neutrino energies can be seen as three distinct groups, with the lower group being the electron neutrinos, the middle group being the electron antineutrinos, and the upper group being the μ and τ neutrinos. The LSBCK neutrino spectra was softer than that of the RMF group over the first 50 ms after bounce, due to the expansion of the PNS radius beyond that of the other models (Figure 5.28) in this epoch. This ceases to be true as the LSBCK PNS radius contracts more steeply than the other models. Further

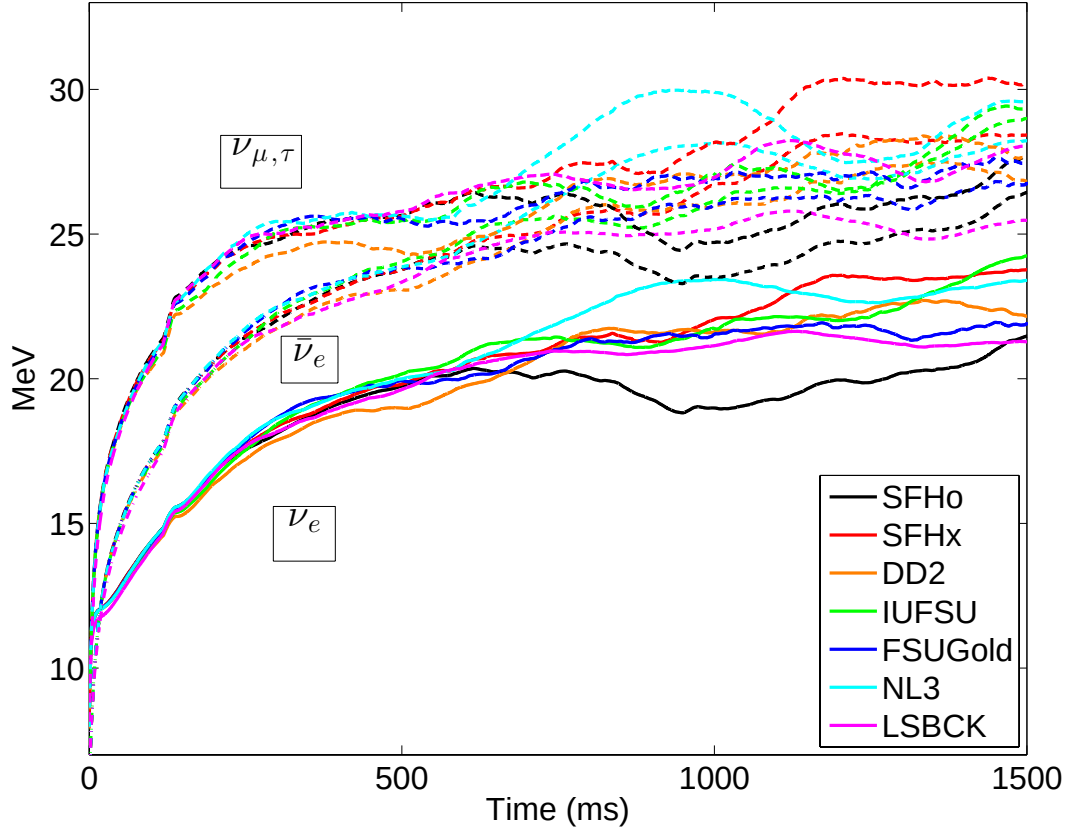


Figure 5.99: Neutrino RMS Energy vs Time, Post-Bounce

as each model compresses and decreases in radius (along with the neutrinospheres), the neutrinosphere temperatures rise and the spectra harden. The neutrino RMS energies can be seen to rise steeply initially, diverge somewhat after 200 ms, and then gradually harden over time except for those of the SFHo model, which begin to decline for each species. Over and above the slowly increasing energies of the non-SFHo models, the NL3 RMS energies for each species harden significantly, showing an upward inflection and leveling off several MeV above the values of the rm energies in the other models.

5.12.2 Neutron Star Characteristics

The PNS mass evolves with the post-bounce accretion rate (Figure 5.10), which tapers to varying degree as the convective engine leads to explosion, with significant dependence on nuclear EOS, as seen in Figure 5.9. The PNS mass increases rapidly during the heavy-accretion early epoch, and begins to climb at a more shallow rate as the accretion rate recedes due to the onset of explosion. The PNS baryon masses (where the PNS is defined as the mass where $\rho > 1 \times 10^{11} \text{ g cm}^{-3}$) at 1.5 sec (Figure 5.9) range from SFHo on the low end of the range at $1.691 M_{\odot}$ to NL3 on the high end of the range at $1.750 M_{\odot}$, but the spread in masses continues to expand. In terms of the PNS radius, the softer EOS of the LSBCK model overshoots the other model in its bounce rebound, and then compress to a smaller, denser PNS, as also evidenced by the steeper slope of the change in central density and the smaller PNS radius (Figure 5.11). The models whose PNS radii are decreasing most slowly likewise have the lowest core densities (Figure 5.100). The stiffer EOS allows for a larger PNS radius for the same PNS mass, and will allow for a larger maximum neutron star mass. The accretion rates seen in Figure 5.10 after 500 ms are reflected in the rate of change of the PNS mass in that time, as the low gain accretion rates of SFHo, SFHx, and FSUGold correlate to flattening PNS mass growth, while the other models continue to grow at a greater rate. These late-low-accretion models have net accretion rates approaching zero (specifically SFHo), which may signal the imminent transition to simultaneous accretion to a neutrino-driven wind. The EOSs with most compact proto-neutron stars at bounce and during the early post-bounce evolution are SFHo and SFHx and remain the most compact models until about 130 ms when the LSBCK model becomes the most compact. The mass/radius relation is plotted in Figure 5.101, showing the evolution of the mass relative to radius during collapse (where it rises as more matter exceeds the PNS threshold $1 \times 10^{11} \text{ g cm}^{-3}$ in density), retraction to bounce, the bounce rebound, and accretion. The central density with respect to PNS mass can be seen in Figure 5.102, where the stiffer EOSs (such as NL3, DD2, and IUFSU) have lower central densities for the same PNS mass as the softer supranuclear EOS (LSBCK, SFHo/x, and FSUGold).

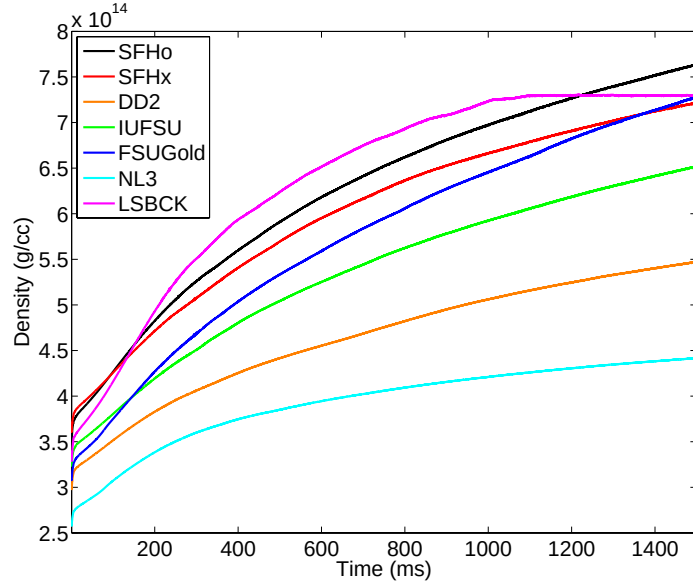


Figure 5.100: PNS Central Density vs. Time

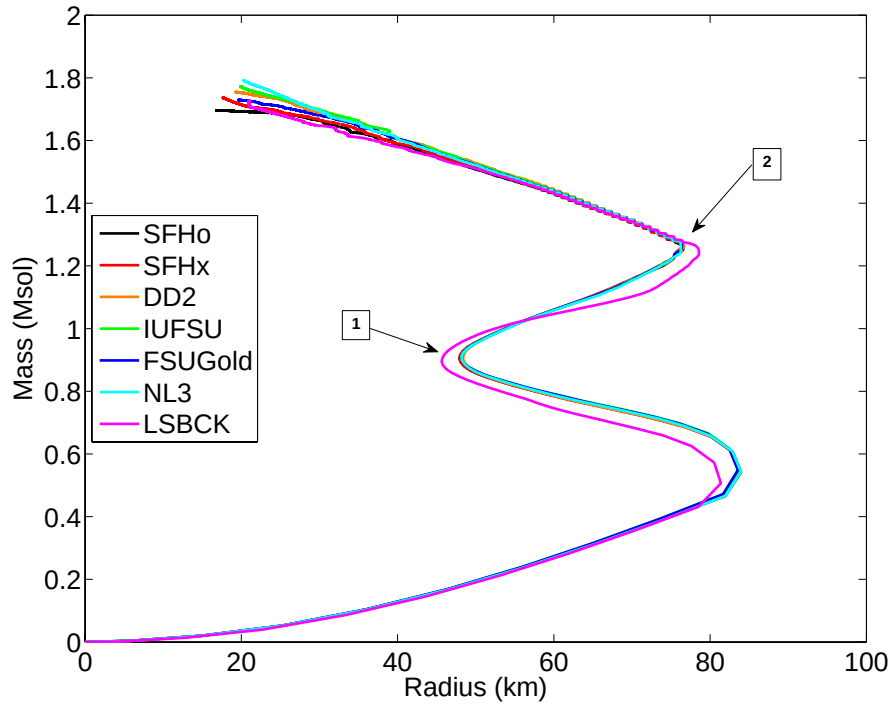


Figure 5.101: Proto-neutron star Mass vs Radius. The point labeled “1” is the bounce event, and the point labeled “2” is the max PNS excursion after bounce, followed by the accretion phase.

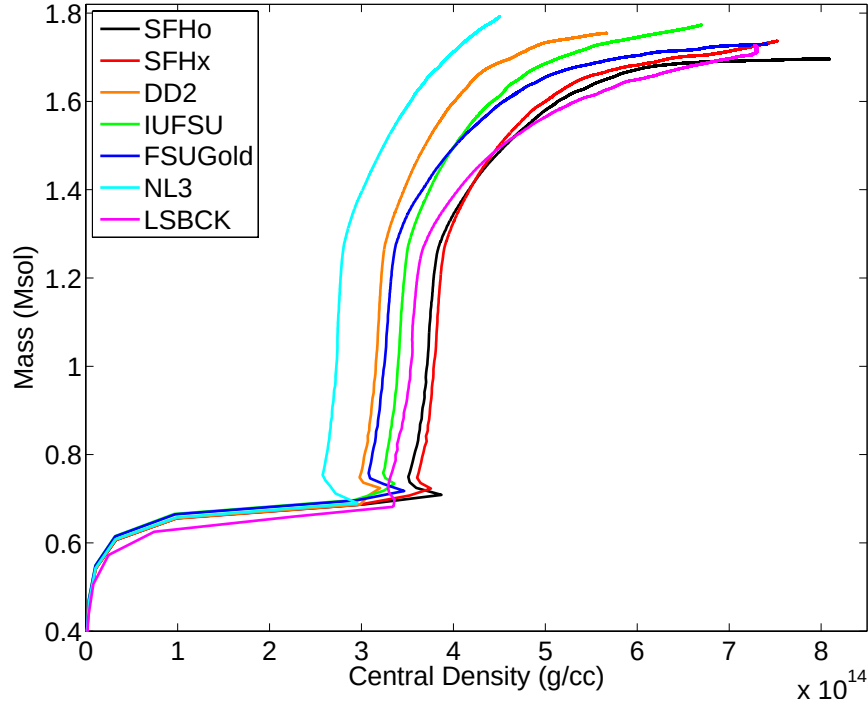


Figure 5.102: Proto-neutron star Mass vs Central Density

5.12.3 Effects of EOS on Growth of Explosion Energy

The shock expansion is fundamentally due to the buildup of hot, neutrino-heated post-shock stellar matter in tandem with the lessening ram pressure from the pre-shock material. Additionally, as the shock propagates to larger radii, the pre-shock ram pressure lessens due to the shallower gravitational potential, as well as the abating accretion rates. Early in the simulations (up to 100 ms after bounce), the shock growth was nearly identical, especially among the RMF group of models. The variation in the shock expansion increased as the convective engine develops. In the mid-range (500 ms) of the of the simulations, the softest equations of state have the highest diagnostic energy (Figure 5.3). In Figure 5.103, the total gain neutrino heating over time shows the EOS models maintaining a similar heating rate during the first 100 ms followed by a differentiation among the models during the luminosity peak (which occurs between 100–200 ms) of $2\text{--}4 \text{ B s}^{-1}$ in heating difference between the lowest peak heating (the FSUGold model) and the highest (the LSBCK model). In light

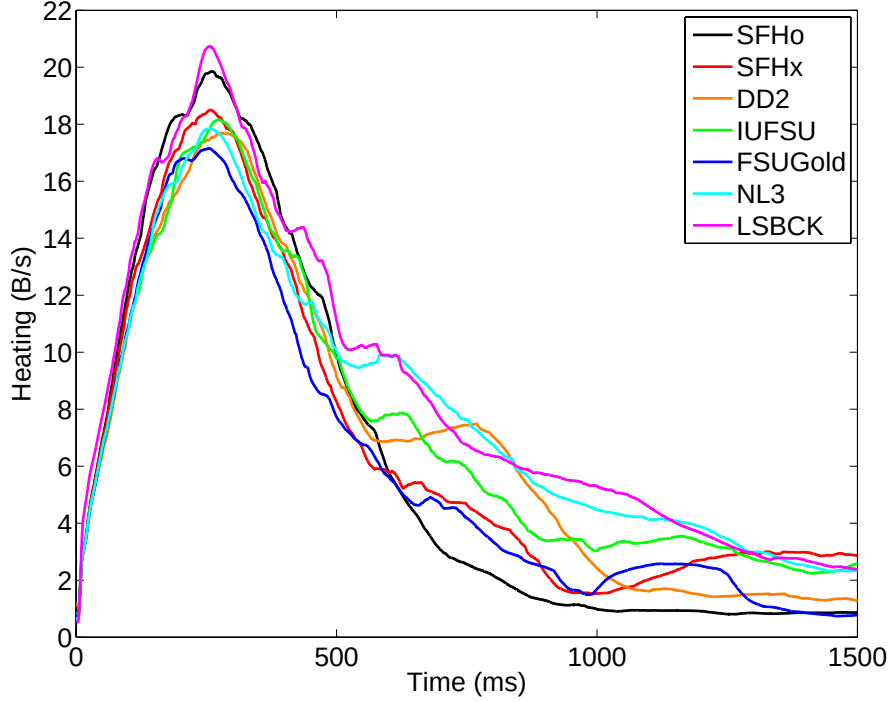


Figure 5.103: Gain Neutrino Heating vs Time, 1500 ms

of the heating-by-weak-process discussion in Section 5.11 the total heating rates seen in Figure 5.103, may be higher for the LSbCK model relative to the RMF models due to the effects of composition on the efficiency of specific weak heating process, both in terms of the peak in ν_e and $\bar{\nu}_e$ heating but also that the LSbCK model tends to have a lower alpha fraction and hence loses less neutrino energy to cooling via the “alpha window”, i.e. the larger alpha mass fraction of the RMF models. The heating rates continue to differentiate and drop with the falling accretion rates and then roughly maintains this stratification until an additional differentiation event occurs near 500 ms, as the heating behavior transitions from being driven by the higher-accretion, mantle-luminosity plus core luminosity early epoch to the lower luminosity, more core-luminosity driven phase, where the luminosity profiles are more strongly characterized by the individual PNS luminosity profiles.

The heating efficiency for the SFHo model declines steeply after 400 ms, reflecting the decline in the gain heating rate, in the same period of time, shown in Figure 5.104. With the gain accretion rate of the SFHo model approaching zero, less and less material in the gain

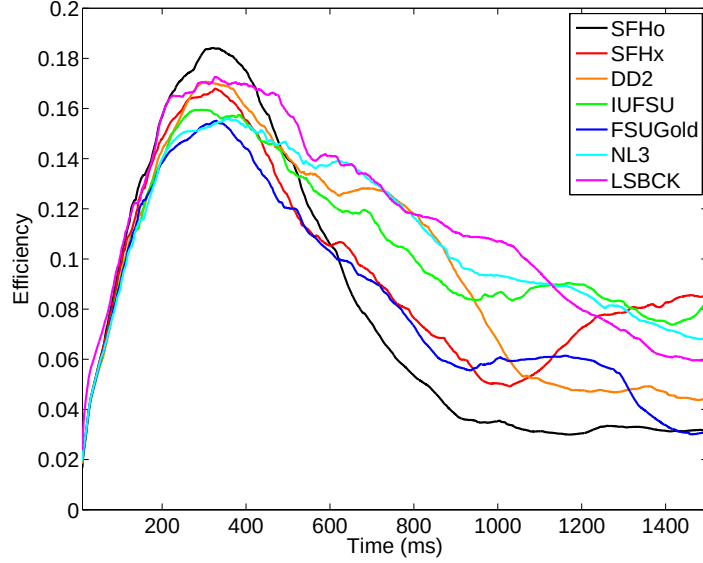


Figure 5.104: Heating Efficiency vs Time

region is available for neutrino-heating. In the same two images, we can see much shallower declines in the heating efficiency for the LSBCK, NL3, and IUFSU models, as well as much slower, linear growth of the heating timescales for those same models, indicating they are still in an earlier phase of explosion.

In section 5.5, the EOS affected the early mass accretion rate, which in turn affected the onset of convection, which by 100 ms, it also affected the occurrence of asphericity. Asphericity affects the formation of accretion streams, which should boost efficiency of heating (as discussed in Section 5.7, which is possibly what aids the heating in the SFHo model in the early part of the simulation. However, the early onset of convection for SFHo seems also to correlate with a less prolate (polar) explosion, its shock front morphology is noticeably the only EOS model that is somewhat spherical at late times, as seen in Figure 5.94, though pending repetition this could simply be chance. The LSBCK model achieves positive total energy by 400 ms (Figure 5.105), and all models have positive total energy by 600 ms.

The spread in luminosity profiles from each EOS model results allows for the description of the different behaviors with respect to explosion among the EOS models simulated in

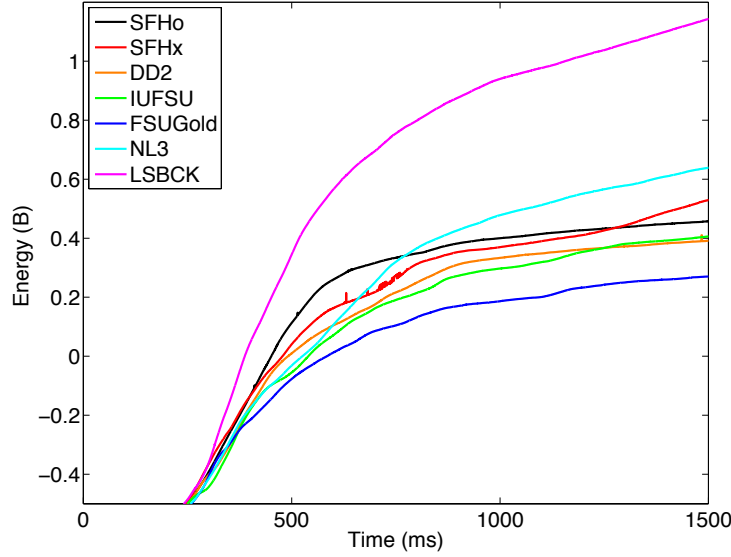


Figure 5.105: Total Energy vs. Time

this study. Firstly, there is the “early time” explosion scenario, exemplified by the SFHo model (and to a lesser extent SFHx), with high early luminosity, early convection onset, high relative PNS radius contraction leading to a compact PNS, with accretion decline, after the luminosity peak, that falls to zero or nearly so, and experiences a late decline in RMS energies as accretion on the PNS ceases to cause additional heating and neutrino spectra hardening. Conversely, the energetics of the NL3 model behave quite differently, with the diagnostic explosion energy peaking late. Lower early luminosity allows for more mid- to later-time accretion that results in lower peak luminosity, but also extends the timescale of the expansion phase. The other RMF models fall somewhere in between these extremes. The LSBCK model, which experiences both high early luminosity from the highest early accretion, low accretion around 300 - 400 ms, and high late luminosity due to both accretion and a luminous PNS, is the outlier and has the most energetic explosion. The unique neutrino luminosity profiles of each of these EOS model groups may likely be a way observational data can help differentiate the most valid equation of state for a given progenitor mass, in terms of the neutrino signatures.

Chapter 6

Conclusion and Outlook

In this dissertation, I first present the new WeakLib EOS framework for the implementation of tabular EOSs into the CHIMERA supernova code. We have improved the treatment of microphysics in the CHIMERA code via the capacity to implement modern microphysical equations of state. With this new capacity, in order to better understand the role nuclear microphysics plays in the evolution of the core-collapse supernova mechanism, I have compared results from one and two-dimensional, self-consistent, numerical simulations of core-collapse supernovae of the $15 M_{\odot}$ Woosley and Heger (2007) progenitor for the first second after bounce and beyond. I have specifically investigated the impact of seven sophisticated microphysical equation of state models on the individual phases of the explosion, from collapse to bounce, rebound, shock stagnation, shock revival, eventual explosion, and late time evolution of the identical stellar progenitors. Extensive differences in the evolution and structure of the proto-neutron star as well as hydrodynamic, compositional, and neutrino emission differences are evident (Figure 5.5, Figure 5.6, Figure 5.99). The EOS impacts the weak interaction physics in each simulation via differences in the evolution of the matter composition, which in turn affects the impact of the each neutrino opacity, influencing the effectiveness of the neutrino-matter heating that leads to explosion in the mantle and expels the infalling layers above. Whereas most previous one-dimensional studies

were unable to investigate the effect of the EOS on explosion differences, by utilizing multi-dimensional simulations, explosion is achieved in all models and hence the effect of the EOS in explosion dynamics can be seen, not possible even in the most modern spherically symmetric simulations.

The complete set of nuclear parameters, i.e. the incompressibility, the symmetry energy, and density dependence slope parameters contribute to the behavior of the EOS in each model and how it influences on the evolution the supernova simulation. The effect of each EOS on the simulation must be described with care, as the classification of EOS as being universally “stiff” or “soft” is not accurate for all EOSs utilized in this dissertation. While a given EOS may have be relatively soft or stiff at a given density, temperature, and electron fraction relative to the others, the relationship relative the others may change in other regions of the density, temperature, and electron fraction phase space. The NL3 EOS model is the exception, having the highest incompressibility K , symmetry energy J , and symmetry energy slope parameter L among all the EOSs (by significant margin for each), as described in Table 3.1. This leads it to be the stiffest of the EOSs employed in this comparison. The LSBCK EOS can be described as the softest, as it has the lowest incompressibility, and nearly the lowest symmetry energy, while the other EOSs fall in between. At supranuclear density, FSUGold is softer than the related IUFSU model which is turn softer than NL3, but at subnuclear densities IUFSU is softer than FSUGold, for example. The LSBCK, SFHo, and SFHx EOSs are softer at high density, as can be seen from the higher PNS contraction rates in Figure 5.11, as well as their higher compactness in Figure 5.13 in their simulations.

The final differences in proto-neutron star mass and radii for the individual models is due to several factors. Firstly, the events at bounce characterize the initial PNS mass. The mass of the proto-neutron star in each model is secondly influenced by the accretion, in that a more luminous model will accelerate the explosion phase and lead to a lower accretion, most notably seen in the SFHo model in Figure 5.9. With the gain accretion dropping below $0.1 M_{\odot} \text{ s}^{-1}$ (shown in Figure 5.10) by 1000 ms, the mass is hardly growing. Additionally, stiffer EOS lead to larger PNS radii, as seen in Figure 5.11, with lower central density, as show in

Figure 5.12. The EOSs with most compact proto-neutron stars at bounce and during the early post-bounce evolution are SFHo and SFHx and remain the most compact models until about 130 ms when the LSBCK model becomes the most compact. This remains the case up until 1000 ms when the PNS in the SFHo and SFHx models become the most compact (Figure 5.13). The IUFSU, FSGold, and NL3 models have larger core masses at bounce, not only due to their higher relative incompressibilities but also due to the influence to their different symmetry energy values.

In this dissertation, I have analyzed the role of the EOS on the luminosity profile of each model (Section 5.12.1) and how the luminosities affect the growth of the explosion. During collapse, LSBCK model experiences the highest deleptonization. This removes electron degeneracy pressure further, leading to faster collapse and higher accretion rate. As I have discussed, this higher contraction rate leads to higher neutrino luminosity of all species, and similar contraction rates are seen in the SFHo and SFHx models. This translates into increased neutrino heating in these models, with the efficient neutrino heating for LS, SFHo, and SFHx in period of high early high accretion. The effectiveness of neutrino heating is determined by the interaction between neutrino opacities and the composition, the later of which is dependent on the choice of equation of state. I have shown that the actual heating is different than the post-processing calculation currently implemented, and that investigation of heating by process will provide more accurate heating rates for a given EOS. I have also shown that the EOS models with lower early luminosity experience more late-time accretion (the exception being the LSBCK model, which is luminous early and experiences significant accretion after 500 ms), and the stiffest EOSs, as in the IUFSU, DD2, and NL3 models, will continue to grow in explosion energy later in the simulation.

The findings presented in this dissertation suggest that the role of the EOS in core-collapse supernova events is complicated. The explosion dynamics and neutrino observables for each EOS model show particular idiosyncrasies for each model, demonstrating how each EOS affects the evolution of the star over the course of the explosion event. The past view that “soft” EOS always produce the most energetic explosions isn’t entirely accurate, as

the FSUGold EOS, one of the softer EOS of the RMF group, produced the least energetic explosion of this study. The findings presented here do indicate that soft EOS models do explode earlier than the stiffer EOS models, and are more energetic early. The role of the EOS in core-collapse supernovae merits further study, especially in three-dimensions.

The primary goal of this modeling effort is to characterize the sensitivity of core-collapse supernova models to the EOS so that these simulations can be used to interpret past observations of such events and predict the neutrino and gravitational waves signature observables of future supernovae and related astrophysical phenomena. The different luminosity and energy profiles described herein can be used to inform neutrino luminosity and energy observations; the neutrino signature of a CCSN would vary in time both in luminosity and energy based on the dynamics of the collapse and explosion. I can continue to take advantage of CHIMERA's stability to investigate how different EOS can influence late time cooling to PNS well after the supernova explosion is successfully underway. As shown in previous studies ([Kuroda et al. 2016](#); [Richers et al. 2017](#)), the gravitational wave signature of core-collapse supernovae is sensitive to the EOS, directly and due to the interplay of the EOS with rotation. Data needed for GW extraction is stored for future analysis. Additionally, the role of progenitor mass on the impact of the EOS remains to be investigated. Also important to investigate will be the role of stochasticity in these simulations so that effects attributed here to the EOS can be correctly understood in context. Lastly, it may prove worthwhile to attempt to separate the mantle accretion luminosity component from the core luminosity component so as to better isolate the effect of the EOS due to PNS dynamics alone.

Bibliography

- Akiyama, S., Wheeler, J. C., Meier, D. L., and Lichtenstadt, I. (2003). The Magnetorotational Instability in Core-Collapse Supernova Explosions. *ApJ*, 584:954–970. [10](#)
- Antoniadis, J., Freire, P. C. C., Wex, N., Tauris, T. M., Lynch, R. S., van Kerkwijk, M. H., Kramer, M., Bassa, C., Dhillon, V. S., Driebe, T., Hessels, J. W. T., Kaspi, V. M., Kondratiev, V. I., Langer, N., Marsh, T. R., McLaughlin, M. A., Pennucci, T. T., Ransom, S. M., Stairs, I. H., van Leeuwen, J., Verbiest, J. P. W., and Whelan, D. G. (2013). A Massive Pulsar in a Compact Relativistic Binary. *Science*, 340:448. [21](#), [24](#), [28](#)
- Baade, W. and Zwicky, F. (1934). On Super-novae. *Proceedings of the National Academy of Science*, 20:254–259. [6](#)
- Baldo, M., Burgio, G. F., and Schulze, H. J. (1998). Onset of hyperon formation in neutron star matter from Brueckner theory. *Phys. Rev. C*, 58:3688–3695. [22](#)
- Baron, E., Cooperstein, J., and Kahana, S. (1985). Type-II supernovae in 12-solar-mass and 15-solar-mass stars The equation of state and general relativity. *Phys. Rev. Lett.*, 55:126–129. [26](#)
- Bethe, H. A. and Wilson, J. R. (1985). Revival of a Stalled Supernova Shock By Neutrino Heating. *ApJ*, 295:14–23. [7](#), [9](#)
- Blondin, J., Mezzacappa, A., and DeMarino, C. (2003). Stability of Standing Accretion Shocks, with an Eye toward Core-Collapse Supernovae. *ApJ*, 584:971–980. [10](#)
- Blondin, J. M. and Mezzacappa, A. (2006). The Spherical Accretion Shock Instability in the Linear Regime. *ApJ*, 642:401–409. [10](#)
- Bruenn, S. W. (1985). Stellar core collapse - Numerical model and infall epoch. *ApJS*, 58:771–841. [47](#), [48](#)
- Bruenn, S. W., De Nisco, K. R., and Mezzacappa, A. (2001). General Relativistic Effects in the Core Collapse Supernova Mechanism. *ApJ*, 560:326–338. [47](#)

- Bruenn, S. W., Lentz, E. J., Hix, W. R., Mezzacappa, A., Harris, J. A., Messer, O. E. B., Endeve, E., Blondin, J. M., Chertkow, M. A., Lingerfelt, E. J., Marronetti, P., and Yakunin, K. N. (2016). The Development of Explosions in Axisymmetric Ab Initio Core-Collapse Supernova Simulations of 12-25 $M_{\odot}$ Stars. *ApJ*, 818:123. [9](#), [54](#), [76](#), [82](#), [99](#), [100](#)
- Bruenn, S. W., Mezzacappa, A., Hix, W. R., Lentz, E. J., Messer, O. E. B., Lingerfelt, E. J., Blondin, J. M., Endeve, E., Marronetti, P., and Yakunin, K. N. (2013). Axisymmetric Ab Initio Core-Collapse Supernova Simulations of 12–25 $M_{\odot}$ Stars. *ApJ*, 767:L6. [9](#), [54](#), [100](#)
- Buras, R., Janka, H.-T., Rampp, M., and Kifonidis, K. (2006). Two-dimensional hydrodynamic core-collapse supernova simulations with spectral neutrino transport. II. Models for different progenitor stars. *A&A*, 457:281–308. [54](#), [82](#)
- Buras, R., Rampp, M., Janka, H.-T., and Kifonidis, K. (2003). Improved Models of Stellar Core Collapse and Still No Explosions: What Is Missing? *Phys. Rev. Lett.*, 90(24):241101. [47](#)
- Burrows, A., Hayes, J., and Fryxell, B. A. (1995). On the Nature of Core Collapse Supernovae Explosions. *ApJ*, 450:830–850. [10](#)
- Burrows, A. and Lattimer, J. M. (1985). The prompt mechanism of Type II supernovae. *ApJ*, 299:L19–L22. [18](#)
- Burrows, A., Livne, E., Dessart, L., Ott, C. D., and Murphy, J. (2006). A New Mechanism for Core-Collapse Supernova Explosions. *ApJ*, 640:878–890. [10](#)
- Burrows, A., Livne, E., Dessart, L., Ott, C. D., and Murphy, J. (2007). Features of the Acoustic Mechanism of Core-Collapse Supernova Explosions. *ApJ*, 655:416–433. [10](#)
- Carbone, A., Colò, G., Bracco, A., Cao, L.-G., Bortignon, P. F., Camera, F., and Wieland, O. (2010). Constraints on the symmetry energy and neutron skins from pygmy resonances in Ni68 and Sn132. *Phys. Rev. C*, 81:041301. [25](#)

- Clark, D. H. and Stephenson, F. R. (1982). The Historical Supernovae. In *NATO Advanced Science Institutes (ASI) Series C*, volume 90, pages 355–370. [4](#)
- Colella, P. and Woodward, P. (1984). The Piecewise Parabolic Method (PPM) for Gas-Dynamical Simulations. *J. Comp. Phys.*, 54:174–201. [46](#)
- Colgate, S. A. and White, R. H. (1966). The Hydrodynamic Behavior of Supernovae Explosions. *ApJ*, 143:626. [8](#), [50](#)
- Daoutidis, I. and Goriely, S. (2011). Impact of the nuclear symmetry energy on the pygmy dipole resonance. *Phys. Rev. C*, 84:027301. [23](#)
- Demorest, P. B., Pennucci, T., Ransom, S. M., Roberts, M. S. E., and Hessels, J. W. T. (2010). A two-solar-mass neutron star measured using Shapiro delay. *Nature*, 467:1081–1083. [21](#), [24](#)
- Fattoyev, F. J., Horowitz, C. J., Piekarewicz, J., and Shen, G. (2010). Relativistic effective interaction for nuclei, giant resonances, and neutron stars. *Phys. Rev. C*, 82:055803. [23](#), [28](#)
- Friedman, B., Pandharipande, V. R., and Usmani, Q. N. (1981). Very hot nuclear matter and pion production in relativistic heavy-ion collisions. 372:483–495. [22](#)
- Gandolfi, S. (2013). The equation of state of neutron star matter and the symmetry energy. In *Journal of Physics Conference Series*, volume 420, page 012150. [22](#)
- Gandolfi, S., Illarionov, A. Y., Fantoni, S., Miller, J. C., Pederiva, F., and Schmidt, K. E. (2010). Microscopic calculation of the equation of state of nuclear matter and neutron star structure. *Mon Not R Astron Soc*, 404:L35–L39. [22](#)
- Hannestad, S. and Raffelt, G. (1998). Supernova Neutrino Opacity from Nucleon-Nucleon Bremsstrahlung and Related Processes. *ApJ*, 507:339–352. [48](#)
- Hawley, J., Blondin, J., Lindahl, G., and Lufkin, E. (2012). VH-1: Multidimensional ideal compressible hydrodynamics code. *Astrophysics Source Code Library*, page 4007. [46](#)

- Hebeler, K., Lattimer, J. M., Pethick, C. J., and Schwenk, A. (2010). Constraints on Neutron Star Radii Based on Chiral Effective Field Theory Interactions. *Phys. Rev. Lett.*, 105(16):161102. [22](#)
- Hempel, M., Fischer, T., Schaffner-Bielich, J., and Liebendörfer, M. (2012). New Equations of State in Simulations of Core-collapse Supernovae. *Astrophys J*, 748:70. [51](#)
- Hempel, M. and Schaffner-Bielich, J. (2010). A statistical model for a complete supernova equation of state. *Nucl. Phys. A*, 837:210–254. [18](#)
- Herant, M., Benz, W., Hix, W. R., Fryer, C. L., and Colgate, S. A. (1994). Inside the supernova: A powerful convective engine. *ApJ*, 435:339–361. [10](#)
- Hewish, A., Bell, S. J., Pilkington, J. D. H., Scott, P. F., and Collins, R. A. (1968). Observation of a Rapidly Pulsating Radio Source. *Nature*, 217:709–713. [6](#)
- Hix, W. R., Messer, O. E. B., Mezzacappa, A., Liebendörfer, M., Sampaio, J. M., Langanke, K., Dean, D. J., and Martinez-Pinedo, G. (2003). Consequences of Nuclear Electron Capture in Core Collapse Supernovae. *Phys. Rev. Lett.*, 91(20):201102. [18](#)
- Horowitz, C. J. (2002). Weak magnetism for antineutrinos in supernovae. *Phys. Rev. D*, 65:43001. [47](#)
- Horowitz, C. J., Kumar, K. S., and Michaels, R. (2014). Electroweak measurements of neutron densities in CREX and PREX at JLab, USA. *European Physical Journal A*, 50:48. [23](#)
- Janka, H.-T. (2001). Conditions for shock revival by neutrino heating in core-collapse supernovae. *A&A*, 368:527–560. [10](#)
- Janka, H.-T. (2012). Explosion Mechanisms of Core-Collapse Supernovae. *Annu. Rev. Nucl. Part. Sci.*, 62:407–451. [9](#), [87](#)
- Janka, H.-T., Langanke, K., Marek, A., Martínez-Pinedo, G., and Müller, B. (2007). Theory of core-collapse supernovae. *Phys. Rep.*, 442:38–74. [10](#)

- Janka, H.-T. and Müller, E. (1996). Neutrino Heating, Convection, and the Mechanism of Type-II Supernova Explosions. *A&A*, 306:167–198. [10](#)
- Kotake, K., Takiwaki, T., Suwa, Y., Iwakami Nakano, W., Kawagoe, S., Masada, Y., and Fujimoto, S.-i. (2012). Multimessengers from Core-Collapse Supernovae: Multidimensionality as a Key to Bridge Theory and Observation. *Advances in Astronomy*, 2012. [51](#)
- Kotake, K., Yamada, S., and Sato, K. (2003). Gravitational radiation from axisymmetric rotational core collapse. *Phys. Rev. D*, 68:044023. [51](#)
- Kuroda, T., Kotake, K., and Takiwaki, T. (2016). A New Gravitational-wave Signature from Standing Accretion Shock Instability in Supernovae. *Astrophys J*, 829:L14. [131](#)
- Lalazissis, G. A., Raman, S., and Ring, P. (1999). Ground-State Properties of Even-Even Nuclei in the Relativistic Mean-Field Theory. *Atomic Data and Nuclear Data Tables*, 71:1. [29](#)
- Langanke, K., Martínez-Pinedo, G., Sampaio, J. M., Dean, D. J., Hix, W. R., Messer, O. E., Mezzacappa, A., Liebendörfer, M., Janka, H.-T., and Rampp, M. (2003). Electron Capture Rates on Nuclei and Implications for Stellar Core Collapse. *Phys. Rev. Lett.*, 90(24):241102–+. [18](#), [48](#)
- Lattimer, J. and Swesty, F. D. (1991). A Generalized Equation of State for Hot, Dense Matter. *Nucl. Phys. A*, 535:331–376. [18](#), [25](#), [26](#)
- Lattimer, J. M., Burrows, A., and Yahil, A. (1985). Type II supernova energetics. *ApJ*, 288:644–652. [18](#)
- Lattimer, J. M. and Lim, Y. (2013). Constraining the Symmetry Parameters of the Nuclear Interaction. *Astrophys J*, 771:51. [20](#), [22](#), [23](#)
- Lattimer, J. M. and Prakash, M. (2006). Equation of state, neutron stars and exotic phases. *Nucl. Phys. A*, 777:479–496. [19](#)

- Lattimer, J. M. and Steiner, A. W. (2014). Constraints on the symmetry energy using the mass-radius relation of neutron stars. *European Physical Journal A*, 50:40. [23](#)
- Marek, A. and Janka, H.-T. (2009). Delayed Neutrino-Driven Supernova Explosions Aided by the Standing Accretion-Shock Instability. *ApJ*, 694:664–696. [10](#), [51](#)
- Messer, O. E. B., Mezzacappa, A., Bruenn, S. W., and Guidry, M. W. (1998). A Comparison of Boltzmann and Multigroup Flux-limited Diffusion Neutrino Transport during the Postbounce Shock Reheating Phase in Core-Collapse Supernovae. *ApJ*, 507:353–360. [10](#)
- Mezzacappa, A., Calder, A. C., Bruenn, S. W., Blondin, J. M., Guidry, M. W., Strayer, M. R., and Umar, A. S. (1998a). An Investigation of Neutrino-Driven Convection and the Core Collapse Supernova Mechanism Using Multigroup Neutrino Transport. *ApJ*, 495:911–926. [10](#)
- Mezzacappa, A., Calder, A. C., Bruenn, S. W., Blondin, J. M., Guidry, M. W., Strayer, M. R., and Umar, A. S. (1998b). The Interplay Between Proto-Neutron Star Convection and Neutrino Transport in Core-Collapse Supernovae. *ApJ*, 493:848–862. [10](#)
- Möller, P., Nix, J. R., and Kratz, K.-L. (1997). Nuclear Properties for Astrophysical and Radioactive-Ion Beam Applications. *At. Data Nuc. Data Tab.*, 66:131–+. [17](#)
- Möller, P., Nix, J. R., Myers, W. D., and Swiatecki, W. J. (1995). Nuclear Ground-State Masses and Deformations. *Atomic Data and Nuclear Data Tables*, 59:185. [17](#)
- Müller, B. and Janka, H.-T. (2014). A New Multi-dimensional General Relativistic Neutrino Hydrodynamics Code for Core-collapse Supernovae. IV. The Neutrino Signal. *ApJ*, 788:82. [54](#), [76](#)
- Nazarewicz, W., Reinhard, P. G., Satuła, W., and Vretenar, D. (2014). Symmetry energy in nuclear density functional theory. *European Physical Journal A*, 50:20. [23](#)
- O’Connor, E. and Ott, C. D. (2013). The Progenitor Dependence of the Pre-explosion Neutrino Emission in Core-collapse Supernovae. *ApJ*, 762:126. [87](#)

- Oertel, M., Hempel, M., Klähn, T., and Typel, S. (2017). Equations of state for supernovae and compact stars. *Reviews of Modern Physics*, 89:015007. [16](#), [17](#), [19](#), [21](#)
- Ohnishi, N., Kotake, K., and Yamada, S. (2006). Numerical Analysis of Standing Accretion Shock Instability with Neutrino Heating in Supernova Cores. *ApJ*, 641:1018–1028. [10](#)
- Oppenheimer, J. R. and Volkoff, G. M. (1939). On Massive Neutron Cores. *Physical Review*, 55:374–381. [24](#)
- Piekarewicz, J. (2007). Validating relativistic models of nuclear structure against theoretical, experimental, and observational constraints. *Phys. Rev. C*, 76:064310. [28](#)
- Piekarewicz, J. (2010). Do we understand the incompressibility of neutron-rich matter? *Journal of Physics G Nuclear Physics*, 37:064038. [25](#)
- Reddy, S., Prakash, M., and Lattimer, J. M. (1998). Neutrino interactions in hot and dense matter. *Phys. Rev. D*, 58:013009. [47](#)
- Richers, S., Ott, C. D., Abdikamalov, E., O’Connor, E., and Sullivan, C. (2017). Equation of state effects on gravitational waves from rotating core collapse. *Phys. Rev. D*, 95(6):063019. [51](#), [131](#)
- Roca-Maza, X. and Piekarewicz, J. (2008). Impact of the symmetry energy on the outer crust of nonaccreting neutron stars. *Phys. Rev. C*, 78:025807. [28](#)
- Roca-Maza, X., Viñas, X., Centelles, M., Ring, P., and Schuck, P. (2011). Relativistic mean-field interaction with density-dependent meson-nucleon vertices based on microscopical calculations. *Phys. Rev. C*, 84:054309. [23](#)
- Scheck, L., Janka, H.-T., Foglizzo, T., and Kifonidis, K. (2008). Multidimensional supernova simulations with approximative neutrino transport. II. Convection and the advective-acoustic cycle in the supernova core. *A&A*, 477:931–952. [87](#)

- Scheidegger, S., Käppeli, R., Whitehouse, S. C., Fischer, T., and Liebendörfer, M. (2010). The influence of model parameters on the prediction of gravitational wave signals from stellar core collapse. *A&A*, 514:A51. [51](#)
- Shen, H., Toki, H., Oyamatsu, K., and Sumiyoshi, K. (1998). Relativistic Equation of State of Nuclear Matter for Supernova and Neutron Star. In *Neutron Stars and Pulsars: Thirty Years after the Discovery*, page 157. [18](#), [26](#), [27](#)
- Steiner, A. W., Hempel, M., and Fischer, T. (2013). Core-collapse Supernova Equations of State Based on Neutron Star Observations. *Astrophys J*, 774:17. [xii](#), [19](#), [25](#), [28](#), [29](#), [51](#)
- Steiner, A. W., Lattimer, J. M., and Brown, E. F. (2010). The Equation of State from Observed Masses and Radii of Neutron Stars. *ApJ*, 722:33–54. [24](#), [28](#)
- Steiner, A. W., Prakash, M., Lattimer, J. M., and Ellis, P. J. (2005). Isospin asymmetry in nuclei and neutron stars. *Phys. Rep.*, 411:325–375. [28](#)
- Stone, J. R. and Reinhard, P. (2007). The Skyrme interaction in finite nuclei and nuclear matter. *Prog. Part. Nuc. Phys.*, 58:587–657. [22](#)
- Sugahara, Y. and Toki, H. (1994). Relativistic mean-field theory for unstable nuclei with non-linear σ and ω terms. 579:557–572. [27](#)
- Suwa, Y., Kotake, K., Takiwaki, T., Whitehouse, S. C., Liebendörfer, M., and Sato, K. (2010). Explosion geometry of a rotating $13 m_{\odot}$ star driven by the sasi-aided neutrino-heating supernova mechanism. *PASJ*, 62:L49–L53. [54](#)
- Suwa, Y., Takiwaki, T., Kotake, K., Fischer, T., Liebendörfer, M., and Sato, K. (2013). On the Importance of the Equation of State for the Neutrino-driven Supernova Explosion Mechanism. *ApJ*, 764:99. [27](#), [51](#), [87](#)
- Suwa, Y., Takiwaki, T., Kotake, K., and Sato, K. (2006). Magnetorotational Collapse of Very Massive Stars: Formation of Jets and Black Holes. In *Origin of Matter and Evolution of Galaxies*, volume 847, pages 476–478. [51](#)

- Suwa, Y., Takiwaki, T., Kotake, K., and Sato, K. (2007). Magnetorotational Collapse of Population III Stars. *Publications of the Astronomical Society of Japan*, 59:771–785. [51](#)
- Suwa, Y., Yamada, S., Takiwaki, T., and Kotake, K. (2016). The Criterion of Supernova Explosion Revisited: The Mass Accretion History. *ApJ*, 816:43. [109](#)
- Symbolist, E. M. D. (1984). Magnetorotational Iron Core Collapse. *ApJ*, 285:729–746. [10](#)
- Tews, I., Krüger, T., Gezerlis, A., Hebeler, K., and Schwenk, A. (2013). Neutron matter with chiral EFT interactions: Perturbative and first QMC calculations. *ArXiv e-prints*, page arXiv:1310.3643. [22](#)
- Tews, I., Lattimer, J. M., Ohnishi, A., and Kolomeitsev, E. E. (2017). Symmetry Parameter Constraints from a Lower Bound on Neutron-matter Energy. *Astrophys J*, 848:105. [22](#), [29](#)
- Thompson, T. A., Burrows, A., and Pinto, P. A. (2003). Shock Breakout in Core-Collapse Supernovae and Its Neutrino Signature. *ApJ*, 592:434–456. [51](#)
- Thompson, T. A., Quataert, E., and Burrows, A. (2005). Viscosity and Rotation in Core-Collapse Supernovae. *ApJ*, 620:861–877. [82](#)
- Todd-Rutel, B. G. and Piekarewicz, J. (2005). Neutron-Rich Nuclei and Neutron Stars: A New Accurately Calibrated Interaction for the Study of Neutron-Rich Matter. *Phys. Rev. Lett.*, 95:122501. [28](#)
- Tolman, R. C. (1939). Static Solutions of Einstein’s Field Equations for Spheres of Fluid. *Physical Review*, 55:364–373. [24](#)
- Tsang, M. B., Zhang, Y., Danielewicz, P., Famiano, M., Li, Z., Lynch, W. G., and Steiner, A. W. (2009). Constraints on the Density Dependence of the Symmetry Energy. *Phys. Rev. Lett.*, 102:122701. [23](#), [25](#)
- Typel, S. (2005). Relativistic model for nuclear matter and atomic nuclei with momentum-dependent self-energies. *Phys. Rev. C*, 71:064301. [27](#)

- Vidaña, I., Polls, A., and Providência, C. (2011). Nuclear symmetry energy and the role of the tensor force. *Phys. Rev. C*, 84:062801. [22](#)
- Weizsäcker, C. F. V. (1935). Zur Theorie der Kernmassen. *Zeitschrift für Physik*, 96:431–458. [17](#)
- Woosley, S. E. and Heger, A. (2007). Nucleosynthesis and remnants in massive stars of solar metallicity. *Phys. Rep.*, 442:269–283. [49](#), [128](#)

Vita

Ryan Elliott Landfield was born in Miami, FL on March 28, 1979 to the parents of Joan and Elliott Landfield. After graduating from St. Edward's High School in 1997, Ryan attended Stanford University in Palo Alto, California, where he majored in Physics. In 2002, he graduated and received his B.S. degree with honors, and again in 2005 with a M.S. in Aeronautical and Astronautical Engineering. After four years working in the private sector as a spacecraft design engineer for Space Systems/Loral, Ryan decided to pursue a Ph.D. in Physics at The University of Tennessee, Knoxville, TN. After initially working with the relativistic heavy-ion group, joined the astrophysics research group, a partnership between UT-Knoxville and Oak Ridge National Laboratory. As part of his research, Ryan utilized high-performance computing to study the topics of computational and nuclear astrophysics in the context of self-consistent core-collapse supernova (CCSN) simulations. He expects to complete his Ph.D. in Physics in Fall, 2018.