Core-Collapse Supernova Explosion Theory

A. Burrows\(^1\)* & D. Vartanyan\(^2\)

\(^1\)Department of Astrophysical Sciences, Princeton University, Princeton, NJ 08544, USA
\(^2\)Department of Astronomy, University of California, Berkeley, CA 94720-3411
*e-mail: burrows@astro.princeton.edu

Most supernova explosions accompany the death of a massive star. These explosions give birth to neutron stars and black holes and eject solar masses of heavy elements. However, determining the mechanism of explosion has been a half-century journey of great complexity. In this paper, we present our perspective of the status of this theoretical quest and the physics and astrophysics upon which its resolution seems to depend. The delayed neutrino-heating mechanism is emerging as a robust solution, but there remain many issues to address, not the least of which involves the chaos of the dynamics, before victory can unambiguously be declared. It is impossible to review in detail all aspects of this multi-faceted, more-than-half-century-long theoretical quest. Rather, we here map out the major ingredients of explosion and the emerging systematics of the observables with progenitor mass, as we currently see them. Our discussion will of necessity be speculative in parts, and many of the ideas may not survive future scrutiny. Some statements may be viewed as informed predictions concerning the numerous observables that rightly exercise astronomers witnessing and diagnosing the supernova Universe. Importantly, the same explosion in the inside, by the same mechanism, can look very different in photons, depending upon the mass and radius of the star upon explosion. A \(10^{51}\)-erg (one “Bethe”) explosion of a red supergiant with a massive hydrogen-rich envelope, a diminished hydrogen envelope, no hydrogen envelope, and, perhaps, no hydrogen envelope or helium shell all look very different, yet might have the same core and explosion evolution.

1 Core-Collapse Supernova Explosions

Stars are born, they live, and they die. Many terminate their thermonuclear lives after billions of years of cooking light elements into heavier elements by ejecting their outer hydrogen-rich envelopes over perhaps hundreds of years. In the process, they give birth to compact white dwarf stars, half as massive as the Sun, but a hundred times smaller. Such dense remnants cool off over billions of years like dying embers plucked from a fire. A subset of these white dwarfs in binary stellar systems will later (perhaps hundreds of millions of years) ignite in spectacular thermonuclear explosions, many of these the so-called Type Ia supernovae used, due to their brightness from across the Universe, to take its measure.

However, some stars, those more massive than \(~8\, M_\odot\), die violently in supernova explosions that inject freshly synthesized elements, generation after generation progressively enriching the interstellar medium with these products of existence. They too leave behind remnants, but neutron
stars and black holes. The former could become radio pulsars, are only the size of a city, and have on average masses 50% again as massive as the Sun. The latter are perhaps a few to ten times more massive than a neutron star, but even more compact and more exotic.

The supernova explosions of these massive stars, the so-called core-collapse supernovae (CCSNe), have been theoretically studied for more than half a century and observationally studied even longer. Yet, the mechanism of their explosion has only recently come into sharp focus. A white dwarf is birthed in these stars as well, but before their outer envelope can be ejected this white dwarf achieves the “Chandrasekhar mass” near \( \sim 1.5 \, M_\odot \). This mass is gravitationally unstable to implosion. After a life of perhaps \( \sim 10-40 \) million years, the dense core of this star implodes within less than a second to neutron-star densities, at which point it rebounds like a spherical piston, generating a shock wave in the outer imploding core. The temperatures and densities achieved lead to the copious generation of neutrinos, so the treatment of the neutrinos and their interaction with dense matter are of great import. This “bounce” shock wave could have been the supernova, but in all credible models this shock wave stalls into accretion, halting its outward progress. This is an unsatisfactory state of affairs — a supernova needs to be launched most of the time to be consistent with observed rates and statistics.

What has emerged recently in the modern era of CCSN theory is that the structure of the progenitor star, turbulence and symmetry-breaking in the core after bounce, and the details of the neutrino-matter interaction are all key and determinative of the outcome of collapse. Spherical simulations seldom lead to explosion. Multi-dimensional turbulent convection in the core necessitates complicated multi-D radiation(neutrino)/hydrodynamic simulation codes, and these are expensive and resource intensive. It is this complexity and the chaos in the core dynamics after implosion that has retarded progress on this multi-physics, multi-dimensional astrophysical problem, until now. In the first era of CCSN simulations, the state-of-the-art was good spherical codes that handled the radiation acceptably. These models rarely, if ever, exploded. Multi-D codes were not yet useful. Then, two-dimensional (axisymmetric) codes arrived, captured some aspects of the overturning convection about which 1D models are mute, but were slow — only a few runs could be accomplished per year. This era was followed by the advent of some 3D capability, and at the same time many 2D runs could be performed to map out some of parameter space and gain intuition concerning the essential physics and behavior. We are now in the era of multiple 3D simulations per year, wherein we can explore core dynamics and explosion in the full 3D of Nature without the fear that a mistake in a single expensive run which could take a year on a supercomputer would set us back. This progress has been enabled by the parallel expansion of computer power over the decades. It is the pivotal role of multi-dimensional turbulence and the breaking of spherical symmetry in the mechanism of explosion itself, coupled with the driving role of neutrino heating, that necessitated the decades-long numerical and scientific quest for the mechanism of core-collapse supernovae. What Nature does effortlessly in a trice has taken humans rather longer to unravel.

However, there are now strong recent indications that the dominant explosion mechanism and rough systematics of the outcomes with progenitor star are indeed yielding to ongoing multi-
pronged international theoretical efforts. Moreover, code comparisons are starting to show general concordance\textsuperscript{2}. Many recent multi-dimensional simulations employing sophisticated physics and algorithms are exploding naturally and without artifice. These include those from our group\textsuperscript{3–10}, using the state-of-the-art code FORNAX\textsuperscript{11}, and those from others\textsuperscript{12–21}. Neutrino heating in the so-called gain region behind a stalled shock, aided by the effects of neutrino-driven turbulence and spherical symmetry breaking, together seem, in broad outline, to be the agents of explosion for the major channel of CCSNe. Other subdominant channels might be thermonuclear (what do the terminal cores of $\sim 8$–$9$ M$_{\odot}$ stars actually do\textsuperscript{22–24}, but see\textsuperscript{25}) or magnetically-driven (so-called “hypernovae”, $\sim 1\%$\textsuperscript{26, 27}; long-soft gamma-ray bursts, $< 0.1\%$). And indeed, there remain for the neutrino mechanism numerous interesting complications concerning nuclear and neutrino physics, the progenitor structures, and numerical challenges to be resolved before this central problem can be retired.

2 How do Core-collapse Supernovae Explode?

It is generally agreed that the stall of the roughly spherical bounce shock wave sets up a quasi-hydrostatic structure interior to it that accretes the matter falling through the shock from the outer core that is still imploding\textsuperscript{28–31}. The rate of accretion ($\dot{M}$) through the stalled shock and onto the inner core is an important evolving quantity that depends essentially upon the density structure of the progenitor’s core just prior to Chandrasekhar instability (see Figure\textsuperscript{1}) and determines much of what follows. The core is so dense and the neutrino particle energies are so high (10’s to 100’s of MeV, million electron volts) that this structure interior to $\sim 10^{11}$ g cm$^{-3}$ is opaque to neutrinos of all species - the structure is a “neutrino star” with “neutrinosphere” radii that depend upon the neutrino species and particle energy, but are initially $\sim 30$-60 kilometers (km) in radius. The bounce shock initially forms in the deeper neutrino-opaque region, but as it emerges quickly (within milliseconds) to larger radii and lower densities a burst of electron neutrinos ($\nu_e$) is generated. It is this burst that saps energy from the shock and leads to its stalling into accretion. A secondary cause of its stalling is the shock dissociation of the infalling nuclei into nucleons. This effect lowers the effective “$\gamma$” of the gas that connects internal thermal energy with pressure by diverting energy into nuclear breakup, thereby channeling less efficiently the gravitational energy otherwise available to provide pressure support for the shock. The stalled shock radius initially hovers around 100-200 km. Just interior to the shock is the semi-(neutrino)transparent “gain region” where the “optical” depth to neutrinos is $\sim 0.1$. This region surrounds the neutrinospheres through which most of the ongoing prodigious neutrino emissions emerge and these bound the inner dense core containing most of the PNS (“proto-neutron star”) mass. This quasi-stable proto-neutron star bounded by the stalled shock fattens by accretion and shrinks by neutrino loss. The neutrino emissions are powered by thermal diffusion from the interior and the gravitational power of accretion. The goal of theory is to determine how the shock is reenergized and launched into explosion, leaving behind the bound neutron star. The explosion is of the mantle of the PNS approximately exterior to the neutrinospheres, and the bound inner material must be left behind.
If there were no ongoing accretion, then neutrino heating in the gain region behind the shock wave would be more than sufficient to power an exploding shock. There would be no tamping accretion ram pressure and neutrino heating by $\nu_e$ and $\bar{\nu}_e$ neutrino absorption on free neutrons and protons in the gain region would easily power a dynamical outflow. This is similar to a thermal wind. However, the accretion ram and neutrino heating are competing to determine instability to explosion, with the added complication that accretion is also powering a changing fraction of the driving neutrino luminosities. The explosion is akin to a bifurcation between quasi-stationary accretion and explosion solutions, with control parameters related to the accretion rate and the neutrino luminosities (for two), but a simple analytic explosion condition in the context of realistic simulations has proven elusive. Hence, detailed simulations are required.

What has emerged is that only those progenitor models with very steep outer density profiles that translate into rapidly decreasing post-bounce accretion rates can explode in spherical symmetry (1D) via the neutrino mechanism. Among the representative progenitor models shown in Figure 1 only the 9 M$_\odot$ star comes close to fitting that description. However, not even it explodes in our 1D simulations. Multi-D seems required. Classically, the 8.8 M$_\odot$ model of Nomoto$^{32}$ explodes spherically, as do a few others with similar very steep outer density profiles$^{3,33}$. However, due to new ideas concerning the character of thermonuclear burning and electron capture in such compact cores, this lowest-mass progenitor region is undergoing a modern reappraisal$^{22–24}$. Note that there are strong arguments in tension with this alternate perspective$^{25}$. The current explosion paradigm for most massive stars is gravitational-energy sourced, neutrino-driven, and turbulence-aided, and we now summarize some of what we have learned concerning the roles of various specific physical effects.

**Efficiency** Since a hot and lepton-rich PNS radiates $\sim 3 \times 10^{53}$ ergs in neutrinos as it transitions into a tightly bound, cold neutron star and supernova explosion energies are “typically” one Bethe, it is often stated that the neutrino mechanism of core-collapse explosion is one of less than 1% tolerances. This is not true. During the 100s of milliseconds to few-second timescales after bounce over which the neutrino heating mechanism operates, the efficiency of energy deposition in the gain region, the fraction of the emitted energy absorbed there, is $\sim 4–10\%$, far higher. Most of the binding energy of the neutron star is radiated over a period of a minute$^{34}$ after the phase during which we think the explosion energy is fully determined.

**Turbulent Convection** Turbulence is fundamentally a multi-dimensional phenomenon and can’t be manifest in spherical (1D) symmetry (and, therefore, in 1D simulations). The turbulence in the gain region interior to the stalled shock is driven predominantly by the neutrino heating itself, which produces a negative entropy gradient unstable to overturn. This is similar to boiling water on a stove, via absorptive heating from below$^{34}$. Figure 2 depicts the inner turbulent convective region early after bounce before explosion, showing accreted matter tracers swirling randomly about the PNS core. A larger neutrino heating rate will increase both the vigor of the turbulence and the entropy of this mantle material. The matter that accretes through the shock on its way inward to the PNS during the pre-explosion phase contains perturbations$^{35–39}$ that arise during pre-collapse
stellar evolution which will seed the convective instability. The larger and more prevalent these seeds the quicker the turbulence grows to saturation and in vigor. One feature of turbulence is turbulent pressure. The addition of this stress to the gas pressure helps push the shock to a larger stalled shock radius. This places matter in more shallow reaches of the gravitational potential well out of which it must climb and helps to overcome the subsequently smaller ram pressure due to infalling matter from the outer core still raining in. The turbulence also forces the accreted matter to execute non-radial trajectories as it settles, increasing the time during which it can absorb neutrino energy before settling on the PNS and, hence, the average entropy that can be achieved in the gain region\textsuperscript{40}. Therefore, through the combined agency of both neutrino heating and neutrino-driven turbulence, the quasi-stationary structure that is the PNS, plus mantle, plus stalled shock wave is more likely to reach a critical condition wherein the steady infalling solution bifurcates into an explosive one. The huge binding energy accumulated in the PNS does not need to be overcome — only its mantle (and with it the rest of the star) needs to be ejected.

Moreover, the turbulent hydrodynamic stress is anisotropic, with its largest component along the radial direction\textsuperscript{40}. Turbulent magnetic stress might also be a factor\textsuperscript{41}. Importantly, turbulence is more effective at using energy to generate stress/pressure than a gas of nucleons, electrons, and photons. As much as \(\sim 30\% \) of the stress behind the stalled shock when the turbulence is fully developed can be in turbulent stress. Hence, partially channelling gravitational energy of infall into turbulence instead of into thermal energy helps support and drive the shock more efficiently\textsuperscript{8}.

**Neutrino-Matter Interactions** The predominant processes by which energy is transferred from the radiated neutrinos to the matter behind the shock in the gain region are electron neutrino absorption on neutrons via \(\nu_e + n \rightarrow e^- + p\), anti-electron neutrino absorption on protons via \(\bar{\nu}_e + p \rightarrow e^+ + n\), and inelastic scattering of neutrinos of all species off of both electrons and nucleons. The two super-allowed charged-current absorption reactions dominate and provide a power approximately equal to the product of the neutrino luminosity and the neutrino optical depth in the gain region. The latter can be \(\sim 10\%\). Therefore, the higher the luminosity and/or absorption optical depth the greater the neutrino power deposition, which can reach levels of many Bethes per second. Upon explosion most of this power goes into work against gravity and only a fraction of the deposited energy is left as the asymptotic blast kinetic energy. This is very qualitatively similar to a thermally-driven wind, for which the energy at infinity scales with the binding energy of the ejecta. Therefore, the stellar binding energy of the ejected mantle might approximately set the scale of the supernova explosion energy. We discuss aspects of this paradigm in §3.

At higher mass densities (\(\rho\)) above \(10^{11} - 10^{12}\) g cm\(^{-3}\), nucleon-nucleon interactions introduce correlations in density and spin. Such non-Poissonian correlations modify the scattering and absorption neutrino-matter interaction rates (generally suppressing them), hence, they affect the emergent neutrino luminosities\textsuperscript{42-47}. This is relevant to the instantaneous power deposition in the gain region, and, hence, the neutrino-driving mechanism itself. These many-body effects increase with density and some of the associated correction factors have been estimated\textsuperscript{48,49} to be of order 10-20\% at \(10^{12}\) g cm\(^{-3}\), near and just interior to the neutrinospheres. However, such
corrections depend upon a detailed and self-consistent treatment of the opacities along with the nuclear equation of state (EOS), and this goal has yet to be achieved. Nevertheless, using scattering suppression factors researchers have shown that these effects can facilitate explosion. They do this by decreasing the opacities, thereby increasing the neutrino loss rates. This leads to a more rapid shrinking of the PNS, which due to consequent compression heats the neutrinosphere regions. This increases the mean energy of the emitted neutrinos. Since the neutrino absorption rates via the charged-current reactions quoted above increase approximately as the square of the neutrino energy and the luminosities themselves are elevated, the neutrino power deposition in the gain region is augmented, thereby facilitating explosion. The effect is not large, but when an explosion is marginal it can be determinative.

During the early collapse phase, increasing densities lead to increasing electron Fermi energies and higher electron capture rates on both free protons and nuclei. Electron capture decreases the electron fraction ($Y_e$, the ratio of the electron density to the proton plus neutron density) of the infalling gas, and this decreases the electron pressure. A decrease in the electron pressure slightly accelerates the infall and the mass accretion rate ($\dot{M}$) versus time. As already stated, $\dot{M}$ after inner core bounce is a key parameter determining, among other things, the accretion ram pressure external to the shock and the accretion component of the neutrino luminosities. Therefore, the rate of capture on infall can affect explosion timing and, perhaps, its viability. The effect is not large, but when things are marginal altering the evolution of $\dot{M}$ can be important. However, the capture rate on the mix of nuclei in the imploding core is not known to better than perhaps a factor of five. Hence, clarifying this important issue remains of interest to modelers.

Finally, the energy transfer to matter via inelastic scattering off electrons and nucleons provides a subdominant component of the driving heating power behind the shock wave. The effect may be only 10%-15%, but, again, when the core teeters on the edge of explosion such effects matter. Neutrino scattering off electrons (akin to Compton scattering, but for neutrinos) results in a large energy transfer, but has a small rate. Energy transfer to the heavier nucleons is small, but the scattering rate is large. The net effect results in comparable matter heating rates for both effects, with a slight advantage to neutrino-nucleon scattering. However, calculating such spectral energy redistribution is numerically difficult, and represents one of the major computational challenges in the field.

**Explosion** The stalled shock radius can be decomposed into spherical harmonics in solid angle. The onset of the explosion of a stalled accretion shock is a monopolar instability in the quasispherical shock. However, approximately when the monopole becomes unstable, the dipole often seems to as well. Therefore, the explosion picks an axis, seemingly at random for a non-rotating progenitor, and the blast has a dipolar structure with a degree of asymmetry that seems to be low for quickly exploding models and larger for those whose explosion is more delayed. Figure depicts an example blast structure manifesting such a dipole. Generally, but not completely reliably, the lower-mass progenitors (such as a 9-M$_\odot$ progenitor) explode earlier and before turbulence is vigorous and the more massive progenitors seem to explode later and after turbulence has achieved
some vigor. Hence, the latter generally, though not every time, explode more asymmetrically, with a larger dipolar component\textsuperscript{57}. The chaos of the turbulence makes the outcome stochastic, so that the direction of explosion is not easily predicted. Importantly, the chaos of the turbulent flow will result in distribution functions of explosion times, directions, explosion energies, explosion morphologies, residual neutron star masses, \textsuperscript{56}Ni yields, general nucleosynthesis, and kick velocities, etc., even for the same star. It is not even known whether those functions are broad or narrow for a given star.

In addition, exploding more along an axis, as depicted in Figure 3, allows the flow external to the shock to wrap around the prevailing axis and accrete along a pinched waist in an equatorial structure. This breaking of symmetry, impossible in spherical symmetry, allows simultaneous accretion and explosion. Whereas a 1D explosion by its nature turns accretion off in all directions, and thereby throttles back the accretion component of the driving neutrino luminosity, in multi-D the accretion component of the luminosity can be maintained. Hence, the breaking of spherical symmetry supports the driving luminosity and facilitates explosion, just as it is getting started. This symmetry breaking is an important aspect of viable explosion models and is impossible in spherical models. Nature unchained to manifest overturning instability leading to turbulence employs this freedom to facilitate explosions that might be thwarted in 1D. Both the turbulent stress and the option of simultaneous accretion in one direction while exploding in another are important features of the CCSN explosion mechanism.

Convection in the progenitor star upon collapse will create perturbations in velocity, density, and entropy that seed overturn and turbulence in the post-shock matter exterior to the inner PNS and the neutrinospheres. The magnitude of these perturbations generally increases with progenitor mass, but their true character is only now being explored in detail. Recently, a number of groups have embarked upon 3D stellar evolution studies during the terminal stages of massive stars\textsuperscript{17,22,24,58,60}. The potential role of aspherical perturbations in the progenitor models in inaugurating and maintaining turbulent convection behind the stalled shock wave is an active area of research\textsuperscript{4,5,17,37,60} and these studies might soon reveal the true nature of accreted asphericities and their spatial distribution. It might also be that low-order modes in the progenitors would naturally result in angular asymmetries in the mass accretion through the shock and provide a path of least resistance that would (however randomly) set the explosion dipole and direction, whatever its magnitude. Such low accretion-rate paths might actually facilitate explosion in circumstances when it would otherwise be problematic.

Another convective phenomenon that can help achieve the critical condition for explosion is proto-neutron-star convection. This is not the neutrino-heating-driven convection in and near the gain region just behind the stalled shock, but overturn driven by electron lepton loss from beneath the neutrinospheres. As electron neutrinos are liberated from the inner PNS mantle around a radius of $\sim 20$ km, the resulting negative $Y_e$ gradient is convectively unstable. This is akin to instabilities in stars due to composition gradients. All proto-neutron stars show this instability, which lasts for the entire duration of PNS evolution and likely continues long after (many seconds to one
minute) the explosion is launched (if it is). PNS convection\cite{10,61,62} accelerates energy loss (particularly via $\nu_\mu$, $\nu_\tau$, $\bar{\nu}_\mu$, and $\bar{\nu}_\tau$ neutrinos) and electron lepton loss in the PNS, thereby accelerating core shrinkage. In a manner similar to the many-body effect, such core shrinkage leads to higher neutrinosphere temperatures and a stronger absorptive coupling to the outer gain region.

We end this section by emphasizing that the most important determinant of explosion, all else being equal, is the mass density structure of the unstable Chandrasekhar core. The density profile translates directly into the mass accretion rate after bounce and this determines both the accretion tamp and the accretion component of the driving neutrino luminosity. Figure\cite{63} provides an example set of density profiles ($\rho(r)$) from 9 $M_\odot$ to 27 $M_\odot$. This set spans most (but not all) massive stars that give birth to CCSNe. There are a few trends in $\rho(r)$ worth noting. First, the lowest-mass massive stars generally have slightly higher central densities and steeper outer profiles and the higher-mass massive stars have lower central densities and significantly shallower outer density profiles. However, the trend in the slope of the outer density profiles is not strictly monotonic with progenitor mass, with some “chaos” in the structures. Ambiguities in the handling of convection, overshoot, doubly-diffusive instabilities, and nuclear rates have led to variations from modeler to modeler in progenitor stellar models up to collapse that have yet to converge. Furthermore, the effects of fully 3D stellar evolution and rotation have not yet been fully assessed. Therefore, the summary behavior depicted in Figure\cite{63} is provisional.

Given these caveats, important general insights are emerging. The first is that the silicon/oxygen shell interface in progenitors (seen for many models in Figure\cite{1}) constitutes an important density jump, which if large enough can kickstart a model into explosion. In many of our models, the shock is “revived” upon encountering this interface\cite{9,64}. The associated abrupt drop in accretion rate and inhibiting ram pressure at the shock upon the accretion of this interface is not immediately followed by a corresponding drop in the driving accretion luminosity. This is due to the time delay between accretion to the shock at 100–200 km and accretion to the inner core where the gravitational energy is converted into useful accretion luminosity. This time delay pushes the structure closer to the critical point for explosion. However, sometimes the density jump is not sufficiently large and its magnitude in theoretical stellar models has not definitively been pinned down.

Second, the very steep density profiles seen for the lower-mass massive stars lead to earlier explosions. However, the associated steeply decreasing rates of accretion also result in less mass in the gain region and a lower optical depth to the emerging neutrino fluxes. The lower absorption depths in the mantle times the lower neutrino luminosities lead to lower driving powers in the exploding mantle. This results in lower explosion energies generically under the neutrino heating paradigm for those stars with steep outer density profiles. Conversely, those stars with shallow density profiles, more often the more massive CCSN progenitors, generally explode later. However, their shallow mass profiles result in more mass in the gain region with a greater optical depth. The larger depth times the larger accretion luminosities lead to greater driving neutrino power deposition. The net effect is often higher asymptotic supernova explosion energies. Hence,
with exceptions, state-of-the-art models suggest that the explosion energy is an increasing function of progenitor mass and the shallowness of the outer density profile of the initial core. In addition, “explodability” does not seem to be a function of “compactness” (a measure of the ratio of progenitor interior mass to radius), with both high and low compactness models exploding. It had been suggested that only low-compactness structures exploded. Not only does this not seem to be true, but it seems that only the higher compactness models can result in explosion energies near the canonical one Bethe. It may be, however, that very high compactness structures have outer mantle binding energies for which the neutrino mechanism can’t provide sufficient driving power. These objects may lead either to weak explosions or fizzes, with many of these leading to black holes (and not neutron stars). In fact, the gravitational binding energy of the mantle of the Chandrasekhar core may set the scale of the explosion energy, and if too high might thwart explosion altogether. This topic deserves much more attention.

3 Supernova Energies

Two-dimensional (axisymmetric) and three-dimensional simulations do not behave the same. The axial constraint and artificial turbulent cascade of the former compromise the interpretation of the results. However, 2D simulations do allow the breaking of important symmetries and overturning motions and are less expensive to perform. Importantly, due to their much lower cost, 2D numerical runs can easily be carried out to many seconds after bounce, something that we and others have found is required for many stars to asymptote to final blast kinetic energies in the context of the neutrino mechanism. So, in order to get a bird’s-eye view of some of the systematic behavior with progenitor mass, we have conducted for this paper a suite of longer-term 2D models using the stellar models of Sukhbold et al. as starting points. For this collection, we have found usually that when a 2D model explodes, its more realistic 3D counterpart does as well, and when it doesn’t neither does the 3D simulation. In our experience, this is usually, but not always, the case, though there is some disagreement on this in the literature. In our recent set of models, it is only the 12-M⊙ and 15-M⊙ models that do not explode. The 2D models generally seem to explode a bit earlier than the 3D models. For instance, the 20-M⊙ and 25-M⊙ stars explode ∼100 ms and ∼50 ms later in 3D. Also, on average the more massive progenitor models explode later. Nevertheless, the shock is (re)launched, if it is, between ∼150 and ∼400 ms after bounce for all these 2D exploding models. This timescale depends, no doubt, upon simulation details (microphysics, resolution, algorithms, etc.), as well as the character of the seed perturbations. For these simulations, we did not impose extra perturbations and left the inauguration of the initial overturning instabilities to numerical noise.

Figure portrays the development of the mean shock radius for all the models of this study. The left-hand panel shows the launch phase, while the right-hand panel provides a later, larger-scale glimpse. The mean shock speeds settle between 10,000 km s⁻¹ and 15,000 km s⁻¹. Table lists the explosion energy, baryonic and gravitational masses, and post-bounce run time. The energies have asymptoted to within a few tens of percent for all models of their final supernova
energies and range from 0.09 to 2.3 Bethes. The 24-M⊙ model is being further scrutinized and is not included here. The higher energies are statistically, but not monotonically, associated with more massive progenitors. The growth of the blast energy is depicted in Figure 5. Those models that asymptote early do so at lower energies. Those models that eventually achieve higher explosion energies not only do so later, but experience deeper negative energies for a longer time before emerging into positive territory. As described in §2, this is what is expected for models with massive (shallow) density mantles, if they explode, and these are generally, though not exclusively, for the most massive progenitors (  > 16M⊙).

The energies shown in Table 1 and Figure 5 include the gravitational, thermal, and kinetic energies, as well as the nuclear reassociation energies, of the ejecta. They also include the outer mantle binding energies of the as yet unshocked material. In this way, all the components of the blast energy are accounted for, except the thermonuclear term. The latter could be as much as ∼10% of the total, and will slightly increase our numbers. However, 0.1 M⊙ of oxygen provides only ∼0.1 Bethe, so it is only for the most explosive progenitors with significant oxygen and carbon shells for which these mostly gravitation-powered supernovae can have an interesting thermonuclear component; this might still be only ∼10%. One would expect that the 56Ni yields would be higher for the more densely mantled stars, so that the thermonuclear energy contribution and 56Ni yield would be correlated with one another and with progenitor mass. Curiously, if the speculations concerning the possible thermonuclear character of the lowest-mass progenitors bear out (though see ), this correlation might be preserved, though for the other end of the massive-star mass distribution (“mass function”). Note that the mass function is weighted towards the lower masses.

Figure 6 superposes the theoretical explosion energies of Table 1 onto a plot of the observationally inferred Type IIp (plateau) supernova energies versus inferred ejecta masses. For our theory numbers, we shift the initial progenitor mass by 1.6 M⊙ to account for an average residual neutron star. In so doing, we do not account for the pre-explosion mass loss of the star, which could be significant. However, the general trend of the inferred energy with a measure of stellar mass is reproduced by the theoretical (black) dots. There is scatter in both the theory and observations, the latter due to systematic uncertainties in the models employed and observational limitations, and the former due to numerical and astrophysical uncertainties. However, natural chaos in the dynamics would naturally lead to a spread in energies (§1 §3), to a degree as yet unknown, even for the same initial stellar structure. We note that there seems to be a larger observational spread in the inferred energies at lower masses. This could reflect natural chaos in the turbulent neutrino mechanism, measurement uncertainties, the effects of unknown rotation, or the possibility that the lowest-mass progenitors explode thermonuclearly just after the onset of a collapse that does not achieve nuclear densities. However, it is too soon to draw any definitive conclusions on this score. Be that as it may, the observed very roughly monotonically increasing trend of explosion energy with mass and the ability of the neutrino mechanism to reproduce the observed range of explosion energies are both encouraging.
Finally, the infalling accretion matter plumes that hit the PNS core generate sound waves that are launched outward. Much of the energy of these sound waves is absorbed behind the shock wave and can modestly contribute to the explosion energy. Such a component is automatically included in our bookkeeping. Though difficult to estimate separately, we don’t envision that acoustic power can contribute more than $\sim 5-10\%$ to the total.

4 Residual Neutron-Star Masses

Figure [7] depicts the evolution of the residual baryon mass of the PNS core for the suite of 2D models investigated here. Such masses flatten early, since the mass accretion rates drop quickly after the explosion commences. The final baryon masses at the last timesteps are given in Table [1] as are the corresponding gravitational masses. The latter include the gravitational binding energy (negative) of the core. These masses range from a low near $\sim 1.2 \, M_\odot$ to a high near $\sim 2.0 \, M_\odot$, nicely spanning the observed range[7]. The neutron star masses we find are closely, but not perfectly, monotonic with progenitor mass and the shallowness of the Chandrasekhar mantle, except for those models that don’t explode. Presumably, these models will eventually collapse to black holes, but on timescales longer than we have simulated.

5 Ejecta Compositions

The issue of the ejecta elemental composition is fundamental to supernova theory. The shallowness of the outer mantle density profile and the associated mass of the inner ejecta are roughly correlated with the yields of oxygen and intermediate-mass (e.g., Ar, Si, Ca) elements. As suggested in §3, such a structure is also likely to explode (if via the neutrino mechanism) with higher energies. Therefore, more of this inner ejecta will be able to achieve the higher temperatures that can transform oxygen and silicon into iron-peak species as well. This includes $^{56}\text{Ni}$. Therefore, one expects that in the context of the neutrino mechanism of explosion $^{56}\text{Ni}$ yields are roughly increasing functions of progenitor mass, with the exceptions to strict monotonicity alluded to previously. Specifically, if a $9\, M_\odot$ star explodes by the neutrino mechanism it cannot have much $^{56}\text{Ni}$ in its ejecta and if a $\sim 16-25 \, M_\odot$ star explodes by the same mechanism the $^{56}\text{Ni}$ yield should be more significant.

All the inner ejecta from the region interior to the stalled shock wave, before and just after explosion, are very neutron-rich ($Y_e \sim 0.1 - 0.2$). As they expand outward, absorption by $\nu_e$ and $\bar{\nu}_e$ neutrinos on balance tends to push the ejecta $Y_e$ upward. If the expansion is fast, then some of the ejecta can freeze out slightly neutron-rich below $Y_e = 0.5$. However, if the expansion is slow, there is plenty of time for some of the debris to become proton-rich ($Y_e > 0.5$). However, generally $Y_e = 0.5$ seems to predominate in the bulk. Therefore, those models that explode early and fast should provide some neutron-rich ejecta, though more of their ejecta could still be proton-rich, while those models that explode later and more slowly (generally, the more massive progenitors)
will be the most proton-rich. This is what we see, where $Y_e$'s from $\sim 0.5$ to as high as $\sim 0.58-0.6$ are found. This might make such supernovae sites for the rp-process and for light p-nuclei (e.g., $^{74}$Se, $^{78}$Kr, and $^{84}$Sr) [65,67]. However, these numbers should be viewed as preliminary, depending as they do on detailed neutrino transport and the complicated trajectory histories of the ejecta parcels. We note that observations of $^{57}$Ni in SN1987A, inferred to be a $\sim 18$ M$_\odot$ progenitor, require that no material with $Y_e$'s lower than 0.497 could have been ejected [78]. Also, none of the ejecta seen in modern simulations can be the site of all the r-process, though the first peak is not excluded. The timescales and $Y_e$'s are not at all conducive.

Furthermore, as stated, inner supernova matter explodes quite aspherically, with bubble, botryoidal, and fractured structures predominating. However, the spatial distribution of $Y_e$ in the ejecta can have a roughly dipolar component, with one hemisphere more proton-rich than its counterpart. Figure 8 depicts a snapshot of a simulation of a 19-M$_\odot$ model. The bluish veil is the shock, while the fractured surface is an isoentropy surface painted by $Y_e$. As seen, there is an orange-purple dichotomy which reflects the fact that the ejecta have a dipole in $Y_e$ that persists. Even an initially uniform ejecta $Y_e$ distribution may be unstable to the establishment of such a dipole. If near and exterior to the $\nu_e$ neutrinosphere at the “surface” of the PNS a perturbation in $Y_e$ arises in a given angular patch of the inner ejecta, that perturbation can grow due the concommitant effect on the absorptive opacity at those angles, which in turn will either suppress or enhance the $\nu_e$ emissions to push the $Y_e$ evolution of that matter parcel in the same direction. The progressive diminution of this absorptive $Y_e$ shift effect with distance can freeze the $Y_e$ perturbation. The upshot is then a crudely dipolar distribution in $Y_e$ that tracks a crudely dipolar angular distribution in the $\nu_e$ and $\bar{\nu}_e$ luminosities and the so-called LESA (“Lepton Emission Sustained Asymmetry”) phenomenon [18,57,79,80]. Whether this dipolar asymmetry in $Y_e$ in the ejecta is a generic outcome remains to be seen.

6 Pulsar Proper Motions - Kicks

The neutron stars born as proto-neutron stars in the supernova cauldron are the source of the radio pulsars known to be darting throughout the galaxy with speeds that average $\sim 350$ km s$^{-1}$ [51,52] and can range up to $\sim 1500$ km s$^{-1}$ [83]. The most natural explanation for these galactic motions is recoils during the supernova explosion directly related to asymmetric matter ejection [64-90] and/or asymmetric neutrino emission. Hence, momentum conservation in the context of at times very aspherical ejection can easily yield the observed speeds. Moreover, it is known that neutrino emissions can have a dipolar component and that the associated net momentum can be large. Neutrinos travel at very, very close to the speed of light and constitute in sum as much as 0.15 M$_\odot$ c$^2$ of mass-energy. Therefore, a mere 1% asymmetry in angle can translate into a kick of $\sim 300$ km s$^{-1}$. However, it is not known how the ejecta and neutrino momentum vectors sum, in particular whether they add or subtract and what the integrated magnitude of the latter is.

Nevertheless, one can speculate about the trends with progenitor star of the magnitude of
the kicks experienced\(^\text{71}\). We have seen that the lowest-mass massive stars tend to explode a bit more spherically, eject less core mass, and emit less energy in neutrinos. The radiated binding energies of the PNS are lower, given the lower accretion rates and lower PNS mass. Hence, we expect the kicks to be smaller for the lower mass progenitors. Conversely, the more massive progenitors tend to explode a bit more aspherically, ejecting more core mass and emitting more mass-energy in neutrinos. Hence, we posit that they produce neutron stars with the greatest kick speeds. There is likely to be some noise in these suggestions, but on average these trends with progenitor mass (actually progenitor structure; see Figure 1) are compelling in the context of the neutrino mechanism of CCSN explosions. Moreover, one would predict that stellar-mass black holes born in the context of core collapse would have low kick speeds, since they are generally expected to have much more inertia/mass than neutron stars and the momentum in any matter ejecta their birth may entail should be smaller. However, the neutrino kicks may be as significant as for neutron-star birth; therefore, the momentum in any such black hole birth kick might be comparable.

7 Pulsar Spins

Stars have angular momentum and spin. As they evolve, the angular momentum is redistributed internally (likely by magnetic torques) and lost in winds. It is not known what the internal birth spin distribution of massive stars is, but crude theoretical calculations suggest that angular momentum is gradually transported out of their cores as they evolve\(^\text{92}\), much of it lost to stellar winds. In stars for which the internal spin rates can be measured due to observed rotational splitting of surface pulsational modes, models of the interior spin evolution leave their interiors rotating ten times too fast\(^\text{93}\). Therefore, the theory of angular momentum redistribution is incomplete. In addition, the cores of massive stars shrink and spin up. So, the spin of a Chandrasekhar core just before collapse is a product of the initial angular momentum distribution, wind angular momentum loss, redistribution torquing during evolution, and progressive evolutionary compression of the core. Furthermore, the spin of the collapsing core can be affected by the stochastic shedding of hydrodynamic waves generated during oxygen and silicon core convection just before collapse\(^\text{94}\). Without many observational constraints, the final core spins before core collapse are unknown.

However, radio pulsars have average surface dipole magnetic fields of \(~10^{12}\) gauss and are observed on average to be rotating slowly, with average periods of \(~500\) milliseconds\(^\text{95,96}\). A neutron star needs to be spinning with a period of \(~5\) milliseconds to have a rotational kinetic energy of \(~10^{51}\) ergs, so periods of \(~500\) ms imply rotational kinetic energies that are four orders of magnitude below supernova energies and these are not dynamically important. Nevertheless, the birth spin of a neutron star is an important observable and predicting this number should be a goal of theory.

Even given the spin rate of the Chandrasekhar white dwarf that collapses, one can’t easily predict the birth spin of the neutron star that eventually emerges. Assuming angular momentum
conservation, collapsing from this initial configuration to a neutron star spins the residue up by approximately a factor of one thousand. So, a ∼50-second progenitor core translates into a ∼50-ms PNS. However, the effects of magnetic torques during collapse and after bounce depend on the unknown core magnetic fields; whatever the initial B-fields, they are radically affected by subsequent compression, rotation, turbulence, and various classes of dynamo action. This is particularly true after bounce, since the violent turbulence in the PNS is likely to radically change both the magnitude and the structure of the magnetic field. Very large fields in the magnetar range ($10^{15}$ gauss) can spin down the nascent PNS on timescales of seconds.

Furthermore, after bounce the accreting PNS can be spun up by accreting matter plumes stochastically, with a jumble of streams with both positive and negative angular momenta that don’t necessarily cancel. This can lead after the mass cut between the nascent PNS and the ejecta to a spinning neutron star, even if the initial star was non-rotating. Figure 9 depicts emergent rotation at later times. The possibility that a previously non-rotating core could be left rotating has been vigorously studied, with a range of spin periods predicted by this process alone from a second or two to tens of milliseconds.

Be that as it may, one would like to determine whether the final neutron star spin is predictable. To date, we don’t know. However, if the total stellar mass and angular momentum loss are determining factors, one would expect neutron stars born in low-metallicity (low abundance of non-H/He elements) massive stars to be faster rotators, all else being equal. Moreover, if angular transfer from the massive star core to its mantle is a continuous process, since the more massive stars evolve more quickly they are likely to leave cores with more angular momentum, again all else being equal. So, lower-metallicity, more-massive massive stars would birth neutron stars with faster spins, but, as implied here, there are still too much uncertainties and too many effectors to draw reliable conclusions.

8 Pulsar B-fields

The prejudice of many researchers is that the frozen-in magnetic flux of the unstable Chandrasekhar core determines the neutron-star fields. This can’t be correct. As stated in §7, turbulence behind the shock and in the PNS itself are natural venues for dynamo growth. At the very least, such violent convective motions will advect and tangle any initial seed fields post-bounce and the multipolarity structure will be radically altered. Even without exponential dynamos, rotation will wind up an initial field and the toroidal and poloidal components will evolve significantly. For large fields near $\sim 10^{15}$ gauss, the field can act back on the rotational profile significantly. Therefore, it is not at all clear what the origin of radio pulsar and magnetar B-fields is, nor what the systematic dependence of these fields might be on progenitor characteristics; clearly, this is a rich and important topic for future research and there have been numerous recent papers attacking aspects of it.

A small subset (∼1%) of core-collapse supernova are so-called “hypernovae” that seem to
be missing links with long-soft gamma-ray bursts (GRBs). Their inferred explosion energies are near \( \sim 10^{52} \) ergs (\( \sim 10 \) Bethes) and seem to be too energetic to be powered mostly by the neutrino mechanism. The best explanation is that these are powered by MHD jets that tap the large spin kinetic energy of fast-spinning (few-millisecond) proto-neutron stars. Such fast rotators may experience strong dynamo action that can achieve magnetar fields. Hence, the natural consequence of rapid rotation may also be large B-fields, that together naturally lead to strong jets that can drive quasi-dipolar outflows. As very tentatively suggested in §7 the progenitors for hypernovae may therefore be the more massive massive stars with low-metallicity, and a subset of them may yield gamma-ray bursts. If the latter is the case, a transition within seconds to a black hole as the maximum mass of a neutron star is reached is indicated, since relativistic jets are the best explanation for such GRBs and only black holes seem able to produce them. Therefore, this class of GRBs would have a non-relativistic jet precursor lasting a few seconds that could be more energetic than the relativistic jet that follows and would eventually overtake the former as it blasted out of the progenitor star. This non-relativistic/energetic to relativistic/less-energetic phasing of jets has yet to be observed in either the hypernova or the GRB context, but is suggested by emerging theory.

In any case, even if the neutrino mechanism described in this paper predominates among CCSNe, the residual neutron star will likely be rotating, however fast, and will have a magnetic field with a dipole contribution. After the neutrino mechanism has cleared out the inner cavity around the nascent spinning/magnetic neutron star, this object will be able to transfer power via weak jets or pulsar action to the inner debris. This may constitute a mere 0.01% to 1% of the total supernova energy. Eventually, its effects will be manifest in the blast remnant, but perhaps as a sub-dominant, under-energetic phenomenon. So, for the canonical case, a standard neutrino-driven explosion could be followed by a weaker magnetically-driven secondary effect most of the time. Signatures of this sort need to be sought, but the supernova remnant Cas A, with its sub-dominant jet/counter-jet structure, may be such a structure. When the initial core is rapidly rotating (we think in a small subset of cases), this magnetically-driven component might overtake in energy the neutrino component. There may also be intermediate cases for which the two mechanisms compete. Such a continuum from weak to strong effects of magnetic fields is an intriguing possibility.

9 Black Hole Formation

If and when the PNS mass exceeds the maximum gravitational mass of a neutron star (with suitable small thermal and compositional corrections), it will collapse to a black hole and continue to accrete. This maximum mass is in the range of 2.1 to 2.4 \( M_{\odot} \) gravitational, about 2.4 to 2.7 \( M_{\odot} \) baryonic, and depends upon the only modestly-constrained nuclear equation of state. How much mass is subsequently accreted depends on how much of the progenitor star is ejected. If none of the star is ejected, and 1) most of the progenitors of such “stellar-mass” black holes are the more massive massive stars with high envelope binding energies (§2) and 2) these have experienced significant pre-collapse wind and/or episodic mass loss, then one would expect the canonical mass
of the product black hole to be $\sim 10-20 \, M_\odot$. This is the helium core mass of those model stars that have very high envelope binding energies exterior to the Chandrasekhar core. However, we had earlier witnessed that stars with initial masses in the $13-15 \, M_\odot$ range didn’t explode. This result could easily be model-dependent and is not the final word on what such stars do. But it is possible that the black-hole outcome is peppered about the massive-star mass function. However, the consensus is that most massive stars with initial masses less than $\sim 20 \, M_\odot$ will lead to neutron stars and most stars with initial masses greater than $\sim 30 \, M_\odot$ should lead to black holes.

If most stellar-mass black holes birthed via collapse have masses in the range of $\sim 10 \, M_\odot$ and neutron stars have a maximum mass a bit above $\sim 2.0 \, M_\odot$, then there would be a “mass gap” between them. Such a gap is suggested by the data, but is not proven. It may be that a shock is relaunched, but has insufficient energy to eject enough of the inner mass and that will then fall back, still launching an explosion wave that unbinds the rest of the stellar envelope. Where this mass cut occurs would determine the birth mass of such a black hole. A “fallback” black hole is a distinct possibility, but may be a small subset of the massive-star mass function. It is likely that most collapses lead to neutron stars, but what the neutron-star/black-hole birth ratio is for the population of massive stars is a subject of much current research.

Finally, there are a few points of principle that need to be articulated. The first is that in the context of the collapse of a Chandrasekhar core, it is impossible to collapse directly to a black hole — there must always be a proto-neutron-star intermediary. This is because the bouncing inner core is out of sonic contact with the outer infalling core. At bounce, the object does not know that it will eventually exceed the maximum mass. This means that even when a black hole is the final outcome, the PNS core will always have a significant neutrino and gravitational-wave signature. A signature of the subsequent dynamical collapse to a black hole will be the abrupt cessation of both signals. Second, given that neutrino energy losses in the range $\sim 0.1-0.4 \, M_\odot c^2$ are “inevitable,” the outer stellar envelope will experience a decrease in the gravitational potential it feels. This will lead to its readjustment on dynamical timescales and the likely the ejection of matter to infinity. Hence, there should always be some sort of explosion, even when a black hole forms. Whether it is such a “potential-shift” explosion, one with significant fallback, or one via a disk jet after the black hole and accretion disk form, it is difficult to imagine a purely quiescent black hole birth.

10 Final Thoughts

As should now be clear, from the vantage of theory, a multitude of effects are of importance in determining the viability, character, and strength of a CCSN explosion. The roles of the initial progenitor structure; multi-dimensional neutrino radiation transport; general relativity; instabilities, turbulence, and chaos; the nuclear interaction and equation of state; neutrino-matter processes and many-body effects; resolution and numerical technique; rotation; and B-fields must all be assessed on the road to a resolution of this complex problem. It is this complexity that has paced progress.
on this multi-physics, multi-dimensional, and multi-decade puzzle. However, modern theory has grappled with all these issues and inputs, with the result that state-of-the-art simulations from many groups evince explosions via the neutrino mechanism with roughly the correct general character and properties. Not all researchers agree on the details. nor do they obtain precisely the same results. Nevertheless, to zeroth-order, the neutrino mechanism seems to work, and one is, therefore, tempted to declare that the overall problem of the mechanism of supernova explosions is solved, the rest being details. However, these details include the credible mapping of progenitor mass and properties to important observables, such as explosion energy, neutron-star mass, nucleosynthesis, morphology, pulsar kicks and spins, and B-field magnitudes and multipolarities. Chaos will complicate all this, as will remaining uncertainties in microphysics and numerics. Yet, despite this, we are confident that core-collapse supernova theory in the year 2020 has reached a milestone, from which it need never look back.
Table 1: **Explosion Energies and Neutron-Star Masses.** The 9-, 10-, and 11-M\(_\odot\) progenitors are from the Sukhbold et al. (2016) suite and were evolved on spherical grids with radial extents of 30,000, 50,000, and 80,000 kilometers (kms), respectively. Progenitors from 12 to 26.99-M\(_\odot\) were inherited from the Sukhbold et al. (2018) suite. The 12-, 13-, and 14-M\(_\odot\) progenitors were evolved on a spherical grid spanning 80,000 km in radius. All other progenitors were evolved on spherical grids spanning 100,000 km in radius. All models were evolved in 2D axisymmetry with 1024 radial cells and 128 (\(\theta\)) angular cells. Thus, there are some small differences in resolution for the lower-mass progenitors, where the progenitor grid is truncated at smaller radii so that the temperature would remain within our equation-of-state table. All models except the 12- and 15-M\(_\odot\) progenitors explode. The Run Time quoted is the time after bounce at nuclear densities.

| Model | Explosion Energy | Run Time | Baryonic Mass | Gravitational Mass |
|-------|------------------|----------|---------------|--------------------|
| [M\(_\odot\)] | [B] | [s] | [M\(_\odot\)] | [M\(_\odot\)] |
| 9     | 0.09             | 2.34     | 1.35          | 1.23               |
| 10    | 0.15             | 3.36     | 1.49          | 1.35               |
| 11    | 0.15             | 3.52     | 1.51          | 1.37               |
| 12    | -0.03            | 2.75     | 1.82          | 1.62               |
| 13    | 0.78             | 4.60     | 1.89          | 1.68               |
| 14    | 0.28             | 4.51     | 1.81          | 1.62               |
| 15    | -0.17            | 1.04     | 1.93          | 1.71               |
| 16    | 0.36             | 4.45     | 1.75          | 1.56               |
| 17    | 1.86             | 4.66     | 2.05          | 1.81               |
| 18    | 1.24             | 4.58     | 1.80          | 1.60               |
| 19    | 0.63             | 4.45     | 1.87          | 1.66               |
| 20    | 1.22             | 4.56     | 2.10          | 1.85               |
| 21    | 1.74             | 3.76     | 2.27          | 1.97               |
| 22    | 0.95             | 4.74     | 2.06          | 1.81               |
| 23    | 0.73             | 4.55     | 2.04          | 1.80               |
| 25    | 1.39             | 3.11     | 2.11          | 1.85               |
| 26    | 2.3              | 4.60     | 2.15          | 1.88               |
| 26.99 | 1.17             | 4.60     | 2.12          | 1.86               |
Figure 1: **Progenitor mass density profiles.** Plotted are the mass density (in g cm$^{-3}$) versus interior mass (in M$_\odot$) profiles of the cores of the progenitor massive stars used as initial conditions for the supernova simulations we highlight in this paper. The associated spherical stellar evolution models were calculated by Sukhbold et al. [63,72] up to the point of core collapse, at which point they were mapped into our supernova code FORNAX [11].
Figure 2: **Inner matter trajectories as the explosion is about to launch.** Shown are the interiors of an explosion only \( \sim 150 \) milliseconds after core bounce (vertical physical scale \( \sim 350 \) kilometers). At this time the shock wave is at \( \sim 150 \) kilometers, just before explosion. The inner ball is the newly-birthed proto-neutron star (PNS) (rendered as an isodensity surface at \( 10^{11} \) g cm\(^{-3} \), colored by \( Y_e \)), surrounded by swirling, turbulent matter, most of which will settle onto the PNS. The trajectories depict the recent 5 milliseconds in the positions of individual accreted matter elements. They are colored by local entropy. The turbulence of this inner region is manifest.
Figure 3: **Early 3D explosion of the core of a 16-M\_\odot star using FORNAX.** Portrayed is a still near 500 milliseconds after core bounce at nuclear densities. The red is a volume rendering of the high-entropy of the ejecta in the neutrino-heated bubbles that constitute the bulk of the volume of the exploding material. The green surface is an isoentropy surface near the leading edge of the blast, the supernova shock wave. Note the asymmetric, though roughly dipolar, character of the explosion and the pinched “wasp-waist” structure of the flow between the lobes. The dot at the center in the newly born neutron star. In this model, as in many others, there is clearly simultaneous accretion at the waist, while there is ejection in the wide-angle lobes. Simultaneous accretion in one sector during concomitant explosion elsewhere maintains the driving neutrino luminosity and is a signature of the useful breaking of spherical symmetry possible in multi-dimensional flow. This contrasts sharply with the artificially enforced situation in 1D/spherical simulations. Simulation performed by the Princeton supernova group\[^6\].
Figure 4: **Mean shock radii of 2D models.** Depicted are the angle-averaged shock radii (in kilometers) of the 2D model suite calculated for this paper versus time (in seconds) after bounce. Most of the models explode, while the 12- and 15-M$_\odot$ progenitor structures do not. The top panel shows the behavior during the first half second after bounce and in the inner 1500 kilometers, with models exploding (when they do) between $t = 0.15$ and $t = 0.4$ seconds. The bottom panel portrays the shock motion on a larger physical scale (15000 kilometers) and to latter times. Many of the models were, in fact, carried to $\sim$4.5 seconds after bounce. The mean shock speeds become rather stable, with values between either 10000 or 15000 km s$^{-1}$ for most of the simulation. The models were conducted on grids from 30,000 to 100,000 kilometers, with the smaller values for the smaller-mass progenitors. See the caption to Table 1 for specifics.
Figure 5: **The evolution of the total explosion energy (in Bethes) with time (in seconds after bounce).** As the figure indicates, many models start bound (negative energies), even though their shocks have been launched. It can take more than one second for some to achieve positive energies, the true signature of an explosion. Moreover, as shown on this figure, it can take $\sim 4-5$ seconds for the supernova energy to asymptote, and some take longer than that. All the more massive exploding models take this longer time, and they generally achieve the highest supernova energies. The lower-mass massive progenitors asymptote earliest at generally, though not universally, lower supernova energies. In addition, though a model might explode late, it can still achieve a higher explosion energy than those that explode early. Hence, the time of explosion is not indicative of its eventual vigor. Note that the 12-M$_\odot$ and 15-M$_\odot$ stars in this investigation do not explode.
Figure 6: **Comparison of theoretical and empirical of explosion energy versus ejecta mass.**

Plotted are the empirically inferred explosion energies versus the inferred ejecta masses, with error bars, for a collection of observed Type IIP (plateau) supernovae. Our theoretical numbers, taken from Table 1, are superposed as black dots. It must be recalled that these are 2D models, and that there are quantitative differences between 2D and 3D simulations. We assume for convenience that the theoretical ejecta masses are the progenitor masses, minus the baryon mass of a putative residual neutron star of $1.6 \, M_\odot$. This ignores any mass loss prior to explosion, surely an incorrect assumption by $\sim 1-3 \, M_\odot$. Nevertheless, the rough correspondence between theory and measurement is quite encouraging. Note that the error bars on the measurements are not firm, and do not include any systematic errors in the light-curve modeling procedures. In any case, the general average trend from low to high explosion energy from lower to higher massive-star progenitor mass reflected in the observations is reproduced rather well by the theory, both quantitatively and qualitatively. Note also that at a given mass there is an inferred measured spread in supernova energies. This may represent a real variation in explosion energy at a given progenitor mass due in part to the natural chaos in turbulent flow. Indeed, it is theoretically expected that Nature would map a given star’s properties to distribution functions in the outcomes and products of its supernova death. The empirical estimates were taken from Morozova et al. [119], Martinez & Bersten [120], Pumo et al. [121][122], and Utrobin & Chugai [123][126].
Figure 7: Theoretical baryon mass (in $M_{\odot}$) of the residual neutron star versus time after bounce (in seconds) for the 2D models of this study. The evolution of the residual neutron star mass is generally rather quick, with the final mass determined to within $\sim 5\%$ generally (though not universally) within $\sim$ one second of bounce. The range of residual masses ranges from $\sim 1.3 M_{\odot}$ to $2.2 M_{\odot}$ for this model set. This is equivalent to a range of neutron-star gravitational masses between $\sim 1.2 M_{\odot}$ and $\sim 2.0 M_{\odot}$, roughly what is empirically seen. Generally, the lower-mass progenitors give birth to lower-mass neutron stars, though this is not rigorously monotonic. Note that the 12-M$_{\odot}$ and 15-M$_{\odot}$ models that don’t explode are still gradually increasing their residual masses by the end of those simulations (see Table 1).
Figure 8: 3D Explosion structure of a representative massive star progenitor model. The associated simulation was performed with a $678 \times 256 \times 512$ grid to render slightly finer details. The snapshot is taken $\sim 800$ milliseconds after core bounce, about 500 milliseconds into the explosion. The blue-gray veil is the shock wave. The colored isosurface is of constant entropy, colored with the electron fraction, $Y_e$. We note that there is a large region of purple (higher $Y_e$, more proton-rich) matter on one side and a largish region of orange-yellow (relatively lower $Y_e$, less proton-rich) matter on the other. This global $Y_e$ asymmetry is created by persistent angular asymmetries in the electron-neutrino and anti-electron-neutrino emission from the core during explosion, which by absorption in the ejecta create this asymmetry in the electron fraction of the ejecta. The latter translates into an asymmetry in the nucleosynthetic element angular distribution. Results derived from a simulation done by the Princeton supernova group.
Figure 9: Similar to Figure 2 but highlighting the generic swirling motions just exterior to the PNS a few hundred milliseconds after explosion. The physical scale top to bottom is \( \sim 200 \text{ km} \). Upon accretion into this inner region, the matter blobs can stream to one side or another of the core before finally settling onto it. This stochastic, almost random, accretion of angular momentum can sum over time to leave a net angular momentum and spin, despite the fact that the original progenitor model was non-rotating[100–102,127].
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Correspondence and requests for materials should be address to A.B.

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