Accretion disks formed from tidally-disrupted companions inside AGB stars

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Abstract

The origin of isolated white dwarfs (WDs) with magnetic fields in excess of \( \sim 1 \) MG has remained a mystery since their initial discovery. Recently, the formation of these high-field magnetic WDs has been observationally linked to binary interactions with low-mass companions (\(< 0.1 \, M_\odot\)) during post-main-sequence evolution. Planetary and M dwarf companions orbiting within several AU of main-sequence stars will become engulfed during the primary’s expansion off the main sequence. Such low-mass companions rapidly in-spiral inside a “common envelope” until they are tidally disrupted near the natal white dwarf core. Formation of an accretion disk from the shredded companion provides a source of turbulence and shear which, in principle, acts to amplify magnetic fields and transport them to the WD surface. However, the disk is initially very cold (\( \sim 10^3 \) K) compared to the hot thermal bath present in the center of the red giant (\( \sim 10^7 \) K). As there is significant shear in the system, entrainment of hot stellar material into the cold disk may lead to vigorous mixing and thermal evaporation before the field can be sufficiently amplified. To study the evolution and lifetime of the tidally-disrupted disks, we perform three-dimensional hydrodynamic simulations of accretion disks formed from 1, 3, 6 and 10 Jupiter mass planets. While entrainment leads to a decrease in disk mass, we find that all disks survive long enough to sufficiently amplify the magnetic field.
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1 Introduction

Close binary interactions are responsible for some of the most energetic events in our universe. While stellar evolution has historically been associated with the study of single stars, binary evolution studies increasingly dominate the literature. The recent discovery of gravitational waves from merging black hole and neutron star systems are several examples of close binary interactions. For the following, we focus one particular type of close binary interaction, namely the common envelope phase.

As stars transition off the main sequence, their radii expand by two-to-three orders of magnitude during the Red Giant Branch (RGB) and Asymptotic Giant Branch (AGB) phases (Iben, 1967; Eggleton, 1971; Iben, 1974; Iben & Renzini, 1984; Kippenhahn & Weigert, 1990). Strong winds driven by radiation pressure acting on dust grains accompany this expansion. At this stage, any companion orbiting within $\sim$20 AU can be expected to strongly interact through (i.) gravitational focusing, shocking and accretion of the wind as it passes by the secondary and (ii.) adjustment of the orbital dynamics due to tidal interactions, mass loss from the system and scattering in multi-body systems (Nordhaus & Spiegel, 2013). In a general sense, the closer the companion, the more energetic the interaction becomes.

For companions initially orbiting $<10$ AU, post-main-sequence expansion will result in a merger as the orbit destabilizes and the companion plunges into the giant. Such common envelope (CE) events are thought to be the primary mechanism for producing short-period binaries in the universe (Paczyński, 1971). Once immersed in a common envelope, dynamical friction leads to rapid inspiral as energy and angular momentum are extracted from the orbit and transferred to the envelope (Ostriker
Typical orbital decay timescales are on the order of months-to-years making this phase of evolution difficult to directly detect (Ivanova et al., 2013). If the companion is of sufficient mass, ejection of the common envelope and emergence in a short-period binary may be possible. If not, the companion is destroyed in the process either through tidal disruption near the giant’s core or through thermal evaporation.

The outcome of CE evolution and the initial post-CE orbital distributions are important as they affect many subfields of astrophysics either directly or indirectly. Stellar population synthesis models often employ ad hoc, constant one- or two-parameter prescriptions for the CE phase based on energy or angular momentum arguments that are independent of mass ratio (see Belczynski et al., 2008, Section 5.4). The output of these codes are used to predict rates for various types of systems (cataclysmic variables, planetary nebulae, binary compact objects, etc.) or to interpret observations of external galaxies. However, such calculations assume spherical symmetry, are tightly coupled to stellar-evolution codes under the assumption of hydrostatic equilibrium (i.e., no velocity field), utilize analytic prescriptions for accretion rates during the CE, and don’t include magnetic or relativistic effects. Unfortunately, there is mounting evidence that accurate outcomes for CE evolution require three-dimensional numerical calculations as hydrodynamic simulations have already shown that many of the listed assumptions built into population synthesis codes are violated (Ricker & Taam, 2012; Passy et al., 2012; Ohlmann et al., 2016).

Common envelopes themselves are also “common” events. Planetary companions to main-sequence stars are plentiful with many predicted to merge during the post-
MS (Nordhaus & Spiegel, 2013). Thus, most stars will evolve through at least one CE phase in their lifetime. Stellar companions such as M dwarfs are known to survive common envelope phases consistent with the current known population of short-period companions to white dwarfs.

Compact objects such as white dwarfs (WD), neutron stars (NS) and stellar-mass black holes (BH) can also become engulfed in giant stars providing avenues to produce all flavors of compact object binaries (i.e. NSNS, WDWD, WDNS, BHBH, NSBH, WDBH). The binary-black hole mergers detected with LIGO are believed to have evolved through a common envelope phase that resulted in the separations necessary to merge via gravitational waves (Belczynski et al., 2016). While we now have multiple data points for binary-black hole systems, the frequency of compact objects mergers depends on which post-CE binaries can merge via gravitational radiation in a Hubble time. Current expected merger rates for binary neutron stars or neutron star-black hole gravitational-wave sources have large error bars. For white dwarfs, the viability of single and double degenerate Type Ia supernovae channels also depends on the distribution of post-CE binaries and their consistency with the observed rates. Lastly, dynamical simulations of common envelopes will produce observational signatures that could be detected by current or future all-sky transient searches (LSST, Pan-STARRS, PTF). In fact, a transient light-curve consistent with some CE phase signatures has tentatively been observed in the galaxy (Ivanova et al., 2013).
1.1 Common Envelope Evolution

Formerly, a common envelope is formed when a component of a contact binary fills its Roche-lobe. This allows material to flow over the outer Lagrange point. Once material escapes this region, it only takes a few orbital time scales for the accreting star to be engulfed in the material from the donor star. This circumbinary material removes energy and angular momentum from the binary system which leads to a rapid decrease in the orbital separation between the core of the giant and the orbiting companion. This process and the associated phenomena are generally known as Common Envelope Evolution (CEE). Common Envelope Evolution is thought to be the main formation channel for many close binary systems and the exotica that may result. Figure 1 shows some of the objects believed to be the result of common envelope evolution.

Paczynski (1976) is credited with the first analysis of the energy and timescales of these type of systems in the context of forming Cepheid variables. The analysis assumes two equal-mass binary components in a low density circumbinary medium that is at rest. The binary components follow a roughly circular orbit with gravitational drag proportional to the ambient density $\rho$, relative velocity, and orbital cross-section. This results in a drag luminosity given by:

$$L_D \sim \frac{d}{dt} \frac{GM^2}{A} \sim \frac{GM^2}{\tau_D A}$$  

(1)

where $A$ is the orbital separation, $M$ is the mass of each stellar component and $\tau_D$ is the orbital decay time scale. With the aforementioned assumptions about the
Figure 1: Ivanova et al 2013 graphic outlining some evolutionary processes where Common Envelope Evolution is important in forming some binary objects. The left column shows double-degenerate mergers and accretion onto a CO WD from a non-degenerate companion, these are two different channels for the creation of SN Ia progenitors. The middle column shows one route by which a binary millisecond pulsar may form. The rightmost column shows a process to produce a double pulsar. Abbreviations: ZAMS-zero age main sequence, RLO-Roche-lobe overflow, CE-common envelope, CO WD-carbon-oxygen white dwarf, He-He star, HMXB-high-mass X-ray binary, LMXB-low-mass X-ray binary, MSP-millisecond pulsar, NS-neutron star, SN-supernova
orbital parameter and gravitational drag, an estimate for the decay timescale is given as:

\[ \tau_D \sim \frac{M}{A^3 \rho} P_{\text{orb}} \sim \frac{\langle \rho \rangle}{\rho} P_{\text{orb}} \tag{2} \]

Where \( \rho \) is the density of the ambient envelope, and \( \langle \rho \rangle \) is the total volume average density including the binary mass. For typical parameters, the decay is rapid, occurring on the order of months to a few years.

1.1.1 Energy Budget

If sufficient orbital energy is transferred to the common envelope, it may become unbound. When this happens before the companion can reach the tidal shredding radius, the formation of a close binary is the result. In this case, the energy formalism that persists from Paczynski assumes the change in orbital energy is used entirely to unbind the common envelope (van den Hauvel 1976; Webbink 1984; Livio and Soker 1988; Iben and Livio 1993):

\[ E_{\text{bind}} = \Delta E_{\text{orb}} = -\frac{Gm_1m_2}{2a_i} + \frac{Gm_{1,e}m_2}{2a_f} \tag{3} \]

In subsequent investigations of common envelope systems that result in close binaries, a more detailed breakdown of the energy is invoked (de Kool 1990; Dewi and Tauris 2000; Dewi and Tauris 2001) such that Eq. (3) is modified by an efficiency parameter, \( \alpha_{\text{CE}} \), that accounts for less than perfect transfer of orbital energy to the envelope. One example of an energy sink in this case is radiative losses from
optically thin regions. Another modification to the formalism is the addition of a second parameter, $\lambda$, that may be has been introduced to account for geometry and internal energy sources. This parameter is highly dependent on the initial separation and stellar model assumed. With these additions, Eq.(3) becomes:

$$\frac{m_1 m_{1,\text{env}}}{\lambda R_1} = \alpha_{CE} \left( -\frac{G m_1 m_2}{2a_i} + \frac{G m_{1,\text{c}} m_2}{2a_f} \right)$$

Eq.(4) seeks to capture the general physics inherent in the common envelope problem. However, it is far too simplistic to emulate the complicated nature of this problem comprehensively. To better understand the dynamics and energy manifestations, theorists look to computational means.

1.1.2 Common Envelope Simulations

Recent simulations have been done to better understand how the energy from the orbital decay manifests itself in the common envelope system. These simulations include smoothed particle hydrodynamics (SPH) and grid-based Eulerian codes to model the initial phases of the interaction. (Ricker Taam (2008, 2012) and Passy et al. (2012), Nandez et al. 2014; Ohlmann et al. 2016; Staff et al. 2016; Kuruwita et al. 2016; Ivanova Nandez 2016; Iaconi et al. 2017, 2018, Chamandy et al. 2018) These simulations often have the issue of the orbital separation decay leveling off at radii too large to explain observations. Other simulations are unable to eject the envelope (Ohlmann et al. 2016; Kuruwita et al. 2016). With these results it is clear that understanding of the processes taking place during CEE is incomplete. However, these failures to reproduce observables can be used to improve our models. Assuming
every process acting in the code is correct, if the envelopes are not unbound it must be that there are untapped energy sources or processes not modeled. Nandez et al. (2015) and Ivanova et al.(2016) suggest recombination as a possible mechanism by which energy is added to the envelope. Whereas Glanz and Perets (2018) assert that the expanding and cooling of the envelope produce radiation pressure that acts on dust grains. The solution to these issues is still debated but the simulations have been invaluable in diagnosing the gaps in the theory.

1.2 High Field Magnetic White Dwarfs

Zeeman spectroscopy and spectropolarimetry reveal a population of isolated, highly-magnetized white dwarfs with surface averaged magnetic field strengths that range from $10^6$ to $10^9$ G (Schmidt et al. 2003). These High Field Magnetic White Dwarfs (HFMWDs) constitute 10% of all isolated white dwarfs (Kepler and others). The remaining 90% of isolated white dwarfs either have detected kilogauss fields, or non-detections with upper limits of $10^4$ to $10^5$ G. Since the first classification of an isolated white dwarf as magnetic in 1970 (Kemp et al. 1970), the origin of these highly magnetized objects has been subject of much debate.

There are currently three different theories regarding formation scenarios for isolated high-field magnetic white dwarfs: (i.) the field is primordial and amplified during the late stages of stellar evolution, (ii.) the field is generated during internal crystallization of the white dwarf, or (iii.) the field results from a close-binary interaction. Before elaborating on these scenarios in more detail, it is important to note a few observational constraints that have been made regarding HFWMDs.
Physically, HFWMDs have been observed to be on average more massive than other isolated white dwarfs. While HFMWDs have a mean mass of $0.784 \pm 0.047 M_\odot$ (M. Silvestri et al. 2007). The majority of white dwarfs are of mass $0.643 \pm 0.136 M_\odot$ (P.-E. Tremblay et al. 2013). HFMWDs have also been observed to be slow rotators having a lower bound for the period of a couple hours and a majority of them on the order of 100 years. Lastly, several HFMWDs have been observed to have very complicated magnetic field geometries (Putney & Jordan 1995; Euchner et al. (2005,2006)).

The most important constraint comes from the Sloan Digital Sky Survey (Gänsicke et al., 2002; Schmidt et al., 2003; Vanlandingham et al., 2005). Remarkably, not a single observed close, detached binary system (in which the primary is a WD and the companion is a low-mass main-sequence star) contains a HFMWD (Liebert et al., 2005; Silvestri et al., 2006). If the magnetic field strengths of white dwarfs were independent of binary interactions, then the observed distribution of isolated WDs should be similar to those in detached binaries. In particular, within 20 pc, there are 109 known WDs (21 of which have a non-degenerate companion). SDSS has identified 149 HFMWDs (none of which has a non-degenerate companion. The probability of obtaining samples at least this different from the same underlying population is $5.7 \times 10^{-10}$, suggesting at the 6-σ level that the two populations are different (Nordhaus et al. 2011). Furthermore, SDSS identified 1253 WD+M-dwarf binaries (none of which are magnetic). This strongly suggests that the presence or absence of binarity is crucial in influencing whether a HFMWD results. On first glance, this may seem to indicate that HFMWDs preferentially form when isolated. However, unless
there is a mechanism by which very distant companions prevent the formation of strong magnetic fields, a more plausible explanation is that highly-magnetized white dwarfs became that way by engulfing (and removing) their companions.

1.2.1 Primordial Field

A primordial origin for the magnetic field has been theorized by Angel et al. (1981). In this framework, the magnetic field flux is conserved during the transition from a low-mass main sequence Ap or Bp star to a white dwarf. Because typical values for the magnetic field strength of a Ap or Bp stars are between $10^3$ and $10^4$ G, increase corresponding to the ratio of volumes would yield field strengths of $10^7 - 10^{10}$ G, consistent with high field magnetic white dwarfs. This theory faces criticism because it does not explain why the average mass of a HFMWD is higher than the majority of isolated white dwarfs. It also does not explain why we do not observe a single HFMWD-M-dwarf binary and conversely why every binary with a white dwarf and a non-degenerate companion does not have a HFMWD.

1.2.2 Crystallization Of The White Dwarf

After a white dwarf forms, it begins to cool and crytalize. During this process it has been proposed that the Rayleigh-Taylor instability insites differential rotation that can generate a magnetic field via an $\alpha - \omega$ dynamo (Isern et al. 2017). The magnetic fields are assumed to be generated in accordance with a scaling relation from Christensen et al. (2010):

$$10^{n}$$
\[
\frac{B^2}{2\mu_0} = c f_\Omega \frac{1}{V} \int_{r_i}^{r_o} \left[ \frac{q_c(r) \lambda(r)}{H(r)} \right] \rho(r)^{1/3} 4\pi r^2 dr,
\]

where the integral is the energy of the convective mantle; \(c\) is an adjustable parameter; \(\mu_0\) is the vacuum permeability; \(f_\Omega \leq 1\) is the ratio of the Ohmic dissipation to the total dissipation; \(q_c\) is the convected energy flux; \(H\) is the scale height; \(V\) is the volume of the convective region; \(r_o\) and \(r_i\) are its outer and inner radii; and \(\lambda\) is the mixing length. Christensen laid out this framework for how turbulent motion can transfer energy to a magnetic field. However, it is contingent on a rotational period that is shorter than the convective turnover timescale (Rossby Number \(>1\)). This condition is necessary for a toroidal field to be wound up before the differential rotation from the Rayleigh-Taylor instability is dissipated. While this is true for rapidly rotating M-dwarfs and planetary cores, HFMWD are slow rotators and their rotation period is much longer than the convective timescale. Therefore, the scaling relation applied to this situation is inappropriate and does not yield accurate results. Furthermore, there is no correlation between the effective temperature and the magnetic field strength. If the magnetic field was the result of crystallization, then the timescales for magnetic field deposition would be on the order of cooling time scales. Therefore it is unlikely that crystallization plays a role in the creation of high field magnetic white dwarfs.

1.2.3 Binary Interactions

With the observation provided from the first paragraphs of section 1.2, it is clear that higher magnetic fields are associated with the absence of a companion. This
makes the common envelope scenario very likely as it has the benefit of removing companions. Current theories utilizing this scenario can be broken down further into WD-WD mergers and Common Envelope Interactions between a post-MS star and a low mass companion.

1.2.4 WD-WD merger

Garcia et al. (2012) proposed that HFMWDs are the result of a WD binary merger. According to SPH simulations, this type of merger would result in a single white dwarf surrounded by a hot corona containing half the mass of the disrupted secondary. Exterior to this is a Keplerian disk of the remaining material from the disrupted secondary. Assuming equipartition between the gas and the magnetic field pressure in the hot corona, the magnetic field strength generated would be of order $10^{10}$ G. Because of the steep temperature gradient, the corona is convective, and the magnetic field will have enough time to set up while the corona cools. Garcia also calculated diffusion time scales for the interior of the white dwarf and the Keplerian disk to be very long, therefore the field should be surface bounded. This theory has been supported by Briggs et al. (2018) using population synthesis codes. However, this theory does not explain the observed slow rotational velocities associated with HFMWDs. In Garcia et al. (2012) they claim magnetic braking could slow the rotation if the magnetic field is perpendicular to the angular momentum vector. The odds that all HFMWDs have magnetic fields perpendicular to the angular momentum is low. Also, as previously stated, not a single HFMWD has an extended non-degenerate companion. It is much more likely for a white dwarf binary
to have extended components, therefore if HFMWDs were the result of WD merger, we would expect to see a HFMWD with an extended companion.

1.2.5 Common Envelope With Low Mass Companion

Nordhaus et al. (2011) theorized that HFMWD are the result of a Common Envelope Evolution between an a low mass extended post-MS star and an M-dwarf companion. The magnetic field could be generated in two channels:

(i) During the interaction, gravitational potential energy would extracted from the binary and injected into the envelope, the differential rotation incited by the inspiraling companion is one method of magnetic field generation. This method was found to be problematic as it does not explain how the field could be anchored to the white dwarf.

(ii) If the orbital energy transferred to the envelope is not enough to gravitationally unbind it, the M-dwarf will inspiral to the tidal shredding radius, a point where the differential potential across the star is greater than its self-gravity. If the companion reaches this radius, it is then gravitationally disrupted and falls rapidly onto the proto-WD core. Because gravity is a central force, angular momentum often conserved. The result is an accretion disk that can generate magnetic fields via a disk dynamo. The magnetic field could then be entrained in the ionized material and deposited onto the surface of the proto-WD core. This theory explains all observations. Because this seems to be the most viable, the magnetic fields generated in this scenario are explored in the next section.
1.3 Magnetic Field Amplification Via Disk Dynamo

Disks formed from this type of tidal shredding event would have a large magnetic Reynolds number. A Reynolds number much greater than one corresponds to flows where the magnetic field lines move with the fluid. Because of this, shear in the disk is provided by the Magneto-rotational Instability (MRI) and is converted via a disk dynamo to magnetic field strength (Balbus and Hawley 1990). As angular momentum is transferred out of the disk, material is accreted towards the WD surface and the strong magnetic field is advected with that flow. Nordhaus calculates the magnetic field strength generated at the surface of the WD to be:

\[
\bar{B}_\phi \sim \left( \frac{\dot{M} \Omega}{\dot{R} H} \right)^{1/2} \approx 160 M G \left( \frac{\eta_{\text{acc}}}{0.1} \right)^{1/2} \left( \frac{\alpha_{\text{as}}}{0.01} \right)^{1/2} \left( \frac{M_c}{10 M_J} \right)^{3/4} \left( \frac{M_{WD}}{0.6 M_\odot} \right)^{1/4} \times \\
\left( \frac{r_c}{r_j} \right)^{-1/4} \left( \frac{H/R}{0.5} \right)^{1/2} \left( \frac{R}{10^9 \text{cm}} \right)^{-3/4} 
\]  

(6)

This equation assumes an equipartition of the gas and magnetic field pressure (Balckman et al. 2000). With Figure 2, it is shown that for companions of mass $1 - 10 M_J$ the resultant magnetic field is on the order of those observed in HFMWDs.
Figure 2: Toroidal magnetic field generated as a result of an alpha disk formed by a tidally shredded companion of mass $M_c$. These lines were calculated with the assumption that $\eta = 0.1$ and $M_{WD} = 0.6M_\odot$. The dotted line represents $\alpha_{ss} = 0.1$ and the solid line represents $\alpha_{ss} = 0.01$. The grey shaded region represents the regime that is tested in this work ($M_c = 1, 3, 6, 10 M_J$), values corresponding to the companion masses that map to the observed magnetic fields of HFMWDs (Nordhaus et al., 2011).

This theory seems promising but it relies on the accretion disk being stable for at least a few rotational timescales to transport the magnetic field through the disk.
and onto the the white dwarf core. To evaluate the stability of such an accretion disk, we turn to simulations.

2 Methods

Simulating these systems computationally could provide information about the validity of magnetic advection through such a disk. Using 3D Magnetohydrodynamical (MHD) codes it is possible to construct an experiment with ideal conditions to test if the thoery is consistent with our physical understanding of accretion disks. In addition to this, computational simulations have the potential to make predictions of observables that can be used to further constrain the mystery at hand. We circumvent the issues highlighted in section 1.1.2 by restricting the simulation time domain to after the tidal shredding event. In this section, we outline the setup and initial conditions of this problem in the 3D MHD code AstroBEAR. First, in order to model the tidally shredded accretion disk system it is imperative to have an accurate stellar profile to serve as the background environment. This is obtained using the stellar evolution code MESA (Paxton 2015). Second, the profiles are then modified in accordance with Ohlmann (2016) to allow for stability in the simulation. Finally, we present the simulation parameters (disk masses, disk geometry, resolution, etc.).

2.1 Extended Stellar Profile

To investigate the accretion disks in the interior of giant stars, it is necessary that the interior of the AGB star be properly modeled. A star leaves the main sequence when
the amount of hydrogen burning in its core is insufficient to provide the radiation pressure necessary to stave off gravitational collapse. When this happens the exterior of the core falls inward causing the pressure and temperature to rise, this allows for hydrogen burning to occur in a shell around the core. Shell burning provides more energy and the outer stellar layers rebound further outward, making them more susceptible to common envelope evolution. We model the complex nature of the post-main sequence lifetime of the star with the 1D stellar evolution code MESA (Paxton et al. 2015). MESA allows the user to model a star through all the phases of it’s evolution. Using a module in the test suite called ‘1MpreMS to WD’ the properties of stars with different masses were modeled in time. Figure 3 shows the classic HR diagram for these stars. Every simulated star starts on the main sequence and evolves through the RG and AGB phases.
Figure 3: Hertzsprung-Russell diagram showing how a star’s luminosity and effective temperature change as they evolve off of the main sequence. Objects in the top right of this plot have a larger radius and therefore are more likely to undergo common envelope evolution.

Because stars with a lower effective temperature give off lower energy photons, for the same luminosity, a cooler star must be much larger. Therefore, stars in the top right corner of the HR diagram have much more extended profiles. Figure 4 shows the radius as a function of time and features two peaks in radius corresponding to the ends of the AGB and RG phases. It is at these points where we extract the profiles as they are the most likely to undergo common envelope evolution.
Figure 4: Stellar radius as a function of time for a $2M_\odot$ star in the post-main-sequence phase. The maximum radii in the RG and AGB phases are marked with red and green crosses respectively.

At each time step MESA outputs the 1D profile for the star. For each of the masses, we plot the profiles corresponding to the maximum radius of the RGB and AGB phases in Figures 5, 6, and 7.
Figure 5: Density Profiles at the maximum radii of AGB and RGB phases for 1, 1.5, and 2.0$M_\odot$ stars.

Figure 6: Pressure profiles at the maximum radii of AGB and RGB phases for 1, 1.5, and 2.0$M_\odot$ stars.
Figure 7: Temperature profiles at the maximum radii of AGB and RGB phases for 1, 1.5, and 2.0\,M_\odot stars.

We choose from these the most extended profile as it is the most likely to fill its Roche-lobe in a binary system. The 2.0\,M_\odot profile serves as the model for which we carry out the subsequent simulations.

### 2.2 Profile Modification

Taking the profiles from MESA and mapping them to the computational grid of our MHD code has proven to be difficult as the innermost region of the star cannot be resolved given computationally limited resolution. Because of the large range of scales inherent to giant stars, Common Envelope simulations often have an issue of resolving the innermost regions of the AGB/RGB profile (e.g., Sandquist et al.)
1998, 2000; Passy et al. 2012; Staff et al. 2016a,b; Iaconi et al. 2017). To solve this issue, the innermost region is integrated and replaced by a point particle. This point particle usually is given the ability to accumulate mass with some prescription (e.g. Bondi, Krumholz). In this work, the point particle used does not accrete and adjacent material simply accumulates around the particle naturally.

Adding a point particle requires a modification of the gravitational acceleration as small separations lead to numerically extraneous accelerations. The usual solution is a softening of the potential at small distances. The gravitational acceleration used by AstroBEAR is smoothed with a cubic spline given by:

$$g_c(r) = Gm_c \begin{cases} \frac{1}{r^2} & r \geq h \\ \frac{y(\frac{12}{3} + y(-48 + y(\frac{192}{5} + y - \frac{32}{3}))) - \frac{2}{30y^2}}{h^2} & h/2 \leq r < h \\ \frac{y(\frac{32}{5} + y^2(-\frac{192}{5} + 32y))}{h^2} & r < h/2 \end{cases}$$

(7)

where $h$ is the smoothing radius, $m_c$ is the core mass, and $y$ is the ratio of radius to smoothing radius. Because the gravitational acceleration was changed, the profile is no longer in hydrostatic equilibrium. Using the method outlined in Ohlmann et al. (2016) the profiles were modified to attain stability. This method is outlined in part.

With the modified gravitational acceleration and hydrostatic equilibrium conditions we have:

$$\frac{dm(r)}{dr} = 4\pi r^2 \rho$$

(8)
\[
\frac{dp(r)}{dr} = -\frac{Gm(r)\rho(r)}{r^2} + \rho(r)g_c(r),
\]  

without accounting for convection and radiative transfer. Solving for the mass enclosed, differentiating and combining the equations we have:

\[
\frac{1}{r^2} \frac{d}{dr} \left( \frac{r^2}{4\pi G \rho} \frac{dp}{dr} + \frac{r^2 g_c(r)}{4\pi G} \right) + \rho = 0
\]  

Assuming a polytrope, we can substitute \( \rho/\rho_0 = \theta^n \) and \( \xi = r/\alpha \). With some algebra we get the modified Lane-Emden equation that can solved numerically:

\[
\frac{1}{\xi} \frac{d}{d\xi} \left( \xi^2 \frac{d\theta}{d\xi} + \xi^2 \frac{g_c(\alpha \xi)}{4\pi G \rho_0 \alpha} \right) + \theta^n = 0
\]

Eq. 11 is very similar to the classical Lane-Emden equations but with an extra term to account for the point particle. In addition to Ohlmann’s method, it is also nessesary to adjust the central mass to account for the constant density still present in the inner most region after the modification. This could cause a few percent difference in the potential. Figures 8 and 9 show the \( 2M_\odot \) stellar profiles from MESA, and the modified profiles to account for finite grid mapping.
With the profiles modified, it is possible to map them stably to a computational grid as long as the smoothing radius is sufficiently resolved.

2.3 AstroBear Simulation Setup

We simulate the accretion disk inside the extended envelope with the Magnetohydrodynamical (MHD) code, AstroBEAR (Cunningham et al. 2009; Carroll-Nellenback et al. 2013). AstroBEAR is an Adaptive Mesh Refinement (AMR) code designed to simulate many different types of physics. With AMR, computational resolution can be adjusted in accordance with necessity. Highly dynamic areas of the grid will be better resolved than those areas that are more relaxed, thus increasing computa-
Figure 9: Density profiles modified in accordance with Ohlmann (2016).

This fully parallelized, multi-dimensional, MHD code was built by the astrophysics group at the University of Rochester.

With AstroBEAR we modeled the system outlined in Nordhaus et al. (2011). We began by mapping the modified MESA profiles from section 2.2 to the grid. Next, the point particle was added with mass corresponding the the integrated missing mass from the modification. Finally, we inserted disks of previously specified companion masses and evolved it for timescales relevant to the problem. This process is outlined in more detail below.
2.3.1 Accretion Disk

Using the companion masses defined in section 1.3 (1, 3, 6, and 10 M\textsubscript{J}), we constructed constant density disks rotating initially at Keplerian velocities around the white dwarf core. The temperature of the disk was set to that of a typical Jupiter sized planet, $\sim 10^3$ K. This is notably lower than the ambient temperature of $10^7$ K that is consistent with stellar cores. The pressure was calculated with the density and temperature assuming $\gamma = 5/3$. The height ($10^9$ cm) and radius ($2.0 \times 10^{10}$ cm) were chosen in accordance with Nordhaus et al. (2011) as those that could generate sufficient magnetic fields.

2.3.2 Global Simulation Parameters

The total simulation time is 20,000s, $\sim 10P_{\text{orb}}$ where $P_{\text{orb}}$ is calculated for the outermost region of the disk. The boxsize we simulate is $10^{11}$ cm and the base resolution is $128^3$ with 3 levels of refinement. To save computational resources all our simulations were done with no AMR. Instead we fixed the grid with the 3 levels of refinement inside a cylindrical container larger than the disk shown in figure 10. This setting is reasonable as the outer most regions of the disk are more dynamically relaxed and should play no important role in the disk stability. We run all of our simulations with 128 cores on Bluehive at the University of Rochester.
Figure 10: Pseudocolor of the density taken from the side of the accretion disk. The black mesh outlines the gridcells. We refine in the inner most cylinder for an effective resolution of $1024^3$. 
3 Results

3.1 Simulation Output

The output of these simulations are frames containing all the important physical information at each time step. It is clear looking at frames below (Figures 11-18) that in all cases the disk persists for the duration of the simulation, supporting the possibility that a magnetic field would have time to anchor to the white dwarf surface. For all visualization and data reduction shown below we use a combination of the python library matplotlib and the 3D visualization software, Visit.
Figure 11: Density and temperature of $1M_j$ disk at 10,000s

Figure 12: Density and temperature of $1M_j$ disk at 20,000s
Figure 13: Density and temperature of $3M_j$ disk at 10,000s

Figure 14: Density and temperature of $3M_j$ disk at 20,000s
Figure 15: Density and temperature of $6M_j$ disk at 10,000s

Figure 16: Density and temperature of $6M_j$ disk at 20,000s
Figure 17: Density and temperature of $10M_j$ disk at 10,000s

Figure 18: Density and temperature of $10M_j$ disk at 20,000s
3.2 Pressure Profile

The pressure profile in time is of interest because of our choice of point particle accretion prescription. We specified no accretion onto the particle, so the disk material just accumulated naturally, building the central pressure, shown in Figure 11. Because the central pressure rises, the disk geometry changes and the accretion rate slows.

![Figure 19: Pressure as a function of radius across the disk. The central potential increases with time.](image)

3.3 Temperature

The disk stability is impressive given the large temperature difference between the accretion disk and the ambient circumstellar medium. Figure 12 illustrates how the ambient material is gradually entrained into the disk as it strips off the outer layers.
Figure 20: Pseudocolor plot of the temperature. The x and y axis are in spatial computational units. The relatively cold disk is slowly evaporated by the ambient medium, entrainment is seen in the structure.

3.4 Disk Geometry

Tracking the disk geometry is important for determining the strength of the magnetic field that can be generated. Figures 21 and 22 show the disk height and disk radius respectively.
Figure 21: Disk height as a function of time for different disk masses. More massive disks result in a taller profile. The low resolution of this figure is the result of the disk height surpassing the high resolution boundary above and below the disk. The height seems to be gradually rising from 10000s onwards.
Figure 22: the outermost disk radius as a function of time. Initially the radius increases with more massive disks expanding to larger radii. After $\sim 5,000$ s the radius begins to gradually decrease for all disk masses.

### 3.5 Accretion Rate

The accretion rate here is defined as the average of all the mass accumulation rates in each gridcell within a radial bin. Because this is defined unphysically, Figure 23 is plotted dimensionlessly, normalized to the initial mass accumulation rate at $r = 0$. It is clear that the innermost gridcells have the largest accretion rate comparatively. However, as the central pressure rises the central accretion rate drops.
Figure 23: Normalized accretion as a function of radius at different times. This plot shows how the accretion rate decreases in the center as the pressure rises. At all other regions of the disk, the accretion rate is fairly constant.

The disk mass is another important parameter for determining the strength of a magnetic field. During the first $\sim 5,000s$ the accretion disk accumulates $0.07M_J$ independent of the disk mass. This increase is due to the initial stellar profile conditions not accounting for the disk in the hydrostatic equilibrium conditions. This mass change is minimal and diminishes with time.
During the first $\sim 5,000$ s every disk accumulates roughly the same mass. After $\sim 5,000$ s, the disk mass steadily decreases.

Figure 24: Relative disk mass as a function of time for different initial disk masses.

4 Discussion and Conclusions

In this work, we presented a brief review of common envelope evolution, highlighted some objects that it produces, and asserted that the population of HFMWDs could be the result of a common envelope evolution between an AGB star and a low mass, M-dwarf companion. To computationally verify this claim we sought to simulate the accretion disk formed from the tidally shredded companion. Simulating this late stage of the evolution circumvents the issues with current common envelope simulations. Using MESA to generate the initial conditions for the ambient density
distribution, we modified the profiles in accordance with Ohlmann (2016) for hydrodynamical stability. With the modified profiles we constructed an AstroBEAR module to simulate an accretion disk inside the stellar profile. Results from these simulations of accretion disks in the interior of giant stars have shown that the disk persists for timescales necessary to deposit a strong magnetic field on the white dwarf core. This supports the theory that HFMWDs are the result of a disk dynamo from a tidally shredded companion. Open questions still exist about the use of an accreting point particle, the resultant magnetic field geometry, and the overall plausibility of accretion disk formation.

Because of our choice of point particle accretion prescription we get that the central pressure rises. Simulating this for longer could be of interest because if the pressure is not relieved through energy outflow, the accretion rate could halt. Turning on magnetic fields in these simulations could incite outflow in the form of jets that should be observationally relevant, as they have a high chance of breaking through the weakly bound stellar envelope, and they could be relevant for shaping the resultant planetary nebula.

4.1 Future Work

To further the investigation of this theory, there are two potential routes for future work. One is to turn on the magnetic fields inside the accretion disk and confirm the resulting magnetic field is of the strength and geometry we observe in the HFMWDs. The other is to simulate the tidal disruption event and verify the formation of the accretion disks modeled in this work. Certainly, more simulations for longer time
with more physics at higher resolution are needed.

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