

# Contents

<b>List of Figures</b>	<b>II</b>
<b>List of Tables</b>	<b>II</b>
<b>1 Introduction</b>	<b>1</b>
<b>2 Thermodynamic Calculation of an EoS</b>	<b>3</b>
2.1 The Ultra-Relativistic Internal Energy . . . . .	3
2.2 The Special-Relativistic Internal Energy . . . . .	4
2.3 The Special-Relativistic Equation of State . . . . .	5
<b>3 Calculating the Mass of a Star with an EoS</b>	<b>9</b>
3.1 Deriving the TOV-Equation . . . . .	9
3.2 Newtonian Limit . . . . .	11
3.3 Mass Bounds . . . . .	12
<b>4 Numerical Solutions</b>	<b>14</b>
4.1 Comparing TOV and LE results with a polytropic EoS . . . . .	14
4.2 Verifying the results . . . . .	16
4.3 Relativistic EoS . . . . .	16
4.4 Zero Values of the TOV and LE equation . . . . .	17
<b>5 To a General Zero Value Theorem</b>	<b>20</b>
5.1 LE Equation . . . . .	20
5.2 TOV Equation . . . . .	22
<b>6 Ideas</b>	<b>25</b>
<b>7 Outlook</b>	<b>26</b>
<b>Appendix</b>	<b>i</b>
A Known Exact Solutions of the LE Equation . . . . .	i
B New Exact LE Series Solution at Index 2 . . . . .	iii
C Exact TOV Solution in Absence of Mass . . . . .	v
D Numerical Optimisations . . . . .	vi
E Nonexistence of Global LE Solutions . . . . .	viii
<b>References</b>	<b>a</b>

## List of Abbreviations

<b>GR</b>	General Relativity . . . . .	1
<b>EoS</b>	Equation of State . . . . .	1
<b>TOV</b>	Tollmann-Oppenheimer-Vollkoff . . . . .	1
<b>LE</b>	Lane-Emden . . . . .	11
<b>ODE</b>	Ordinary Differential Equation . . . . .	1

## List of Symbols

$\mathcal{F}$	Free Energy
$\mathcal{U}$	Internal Energy
$\mathcal{Z}$	Partition Function
$\mathcal{S}$	Entropy
$m$	Mass
$p$	Pressure
$\rho$	Density

## Units

This thesis uses the geometrized unit system  $G = c = 1$ . If not otherwise mentioned any variable named  $G$  or  $c$  can be set to 1 and is only written for clarification.

## List of Figures

1	Relativistic Equation of State . . . . .	8
2	Graph of exact Lane Emden Solutions . . . . .	11
3	Comparison of the TOV and LE equation . . . . .	15
4	Validation of numerical LE results . . . . .	16
5	Comparison of TOV polytropic and relativistic EOS . . . . .	17
6	Zero Values of TOV and LE equation . . . . .	19
7	LE Solution for $n = 2$ . . . . .	v

## List of Tables

1	Lane Emden exact Solutions . . . . .	11
2	Numerical Parameters for TOV and LE equation . . . . .	15

# 1 Introduction

The theory of General Relativity (GR) has allowed scientists deep insights into space and time itself. From its emergence in 1916 [Ein16] to its current state it is regarded as one of the most elegant and well tested theories that Physics has to offer [HK72; Ves+80; Bat+09]. Most objects studied in this framework are typically very heavy or close to the speed of light and thus describes the large scale structure of our observable universe. GR is formulated in a geometric context by regarding space and time itself as a Lorentz-Manifold called spacetime with a non-semi-definite scalar product called the metric. This small difference has non-trivial implications when comparing Riemannian and Lorentzian geometry although we will not make any further effort to properly explain them in detail. A curvature can be introduced that depends on mass and energy present at the respective points in spacetime, and the Einstein equations [Ein15] complete the theoretic description.

$$G_{\mu\nu} + \Lambda g_{\mu\nu} = R_{\mu\nu} + \left(\frac{1}{2}R + \Lambda\right) g_{\mu\nu} = 8\pi T_{\mu\nu} \quad (1.0.1)$$

The left hand side of equation 1.0.1 contains the Einstein Tensor  $G_{\mu\nu} = R_{\mu\nu} + 1/2Rg_{\mu\nu}$  which is defined in terms of the Ricci  $R_{\mu\nu}$  and metric tensor  $g_{\mu\nu}$  and the Ricci scalar  $R$ . These quantities carry information about the curvature of space which can be completely described by the Riemann tensor  $R_{\mu\nu\sigma}^{\lambda}$ . The Energy Momentum tensor on the right-hand side contains information about mass and energy present at different points of space and time. With the Levi-Civita connection on spacetime we can write the Riemann tensor as derivatives of the metric tensor which means the left-hand side of the Einstein equations 1.0.1 contains non-linear derivatives of  $g_{\mu\nu}$ .

$$\Gamma_{\mu\nu}^{\lambda} = \frac{1}{2}g^{\lambda\sigma} (\partial_{\nu}g_{\sigma\mu} + \partial_{\mu}g_{\sigma\nu} - \partial_{\sigma}g_{\mu\nu}) \quad (1.0.2)$$

$$R_{\mu\nu\lambda}^{\sigma} = \partial_{\nu}\Gamma_{\mu\lambda}^{\sigma} - \partial_{\mu}\Gamma_{\nu\lambda}^{\sigma} + (\Gamma_{\mu\lambda}^{\alpha}\Gamma_{\nu\alpha}^{\sigma} - \Gamma_{\nu\lambda}^{\alpha}\Gamma_{\mu\alpha}^{\sigma}) \quad (1.0.3)$$

This non-linearity is one of the key distinctions to other fundamental theories such as a Quantum Theory of Fields and is one of the reasons why few exact solutions of Einsteins equations exist. In addition, the derivatives of  $g_{\mu\nu}$  are also partial which means this theory in general has to deal with non-linear partial differential equations. This behaviour can in general not be altered by simply changing variables. While these statements might sound demotivating, a lot of methods to linearize or obtain Ordinary Differential Equations (ODEs) have been developed. In particular this thesis will use suitable chosen assumptions to derive a famous ODE, the Tollmann-Oppenheimer-Volkoff (TOV) equation and analyse its properties.

Of the plentiful present problems one in particular of significant importance has been to describe the macroscopic properties of stellar objects which in practice are often spherically symmetrical. Stars, black holes, galaxies and galaxy clusters are just a few examples of this category. The TOV equation relates pressure, density and mass inside a spherically symmetric object and given an Equation of State (EoS) one can solve the differential equation and thus obtain parameters such as the mass and total radius of said stellar object. The outside solution of the TOV equation is characterised by the well known Schwarzschild solution [Sch16] which is also applied in the context of black holes.

As was explained just before, an additional piece of information, the EoS needs to added

to the TOV equation in order to obtain solutions. Together with non-negative initial values, we specifically prove, that solutions in general exist and are unique. From a Physics perspective, it is clear that thermodynamics has to play a role in the description of stars since one has to deal with large numbers of particles well in equilibrium. Thermodynamics is a statistical theory of nature and describes macroscopic phenomena by knowing microscopic behaviour of particles. Over the years, different ensembles have emerged and can equivalently, although sometimes in practice more easily done by a specific choice, describe the same physical system. The canonical ensemble is characterised by its partition function given by

$$\mathcal{Z}(T, V, N) = \int \exp\left(-\frac{H(x_1, \dots, p_N)}{k_B T}\right) \frac{dx_1 \dots dp_N}{N! h^{3N}}. \quad (1.0.4)$$

The partition function  $\mathcal{Z}$  can then be used to calculate properties such as pressure  $p$ , internal energy  $\mathcal{U}$  and thus also the energy density  $\mathcal{U}/V$ . Together with an adiabatic condition one would hope to be able to derive a EoS. In many examples, a polytropic EoS between pressure  $p$  and density  $\rho$  is assumed

$$\rho = A p^{1/\gamma} \quad (1.0.5)$$

where  $A$  is a constant and the exponent  $\gamma$  can be related to the adiabatic index of a thermodynamic process. This equation is derived by using the ideal gas equation  $p v = N k_B T$  and adiabaticity. While there exist more recent examples of more complex EoS models [HM88], this equation already covers a wide range of examples. The index  $\gamma$  is typically defined by the adiabatic condition  $p V^\gamma = \text{const}$  and thus directly related to the previous assumption. Another important relation is between  $\gamma$  and degrees of freedom. For an ideal gas, the index can be given by  $\gamma = 1 + 2/f$ , where  $f$  is the degrees of freedom. This quantity directly yields information about the underlying microscopic behaviour of the gas. A polytropic EoS describes many systems and is a good enough assumption in many cases. Part of this thesis will be concerned with deriving a fully special relativistic EoS and applying it in the context of the TOV equation.

## 2 Thermodynamic Calculation of an EoS

This chapter aims at developing a fully special relativistic EoS of a non-interacting gas. We briefly summarise important concepts necessary for the derivation. In the canonical ensemble, for which an introduction can be found in [Fli18], the internal energy  $\mathcal{U}$  is obtained by the relation  $\mathcal{U} = \mathcal{F} - TS$  where the quantities  $\mathcal{F}$  and  $S$  can be derived by means of the partition function  $\mathcal{Z}$  while  $T$  is a variable. Microscopically, the partition function is given by the behaviour of the  $N$  particles determined by the Hamiltonian  $H$ . Concepts and definitions of Hamilton Mechanics can be found in [Esc11; Fli20; Spi10]. In general we assume  $H : \mathbb{R}^{3N} \times \mathbb{R}^{3N} \rightarrow \mathbb{R}$  to be a positive smooth function that assigns an energy to the positions and momenta of  $N$  particles. This explicitly takes form in the well known equations

$$\mathcal{Z}(T, V, N) = \int_{\mathcal{V} \times \mathbb{R}^N} \exp\left(-\frac{H(x_1, \dots, x_N, p_1, \dots, p_N)}{k_B T}\right) \frac{dx_1 dp_1 \dots dx_N dp_N}{N! h^{3N}} \quad (2.0.1)$$

$$\mathcal{F}(T, V, N) = -k_B T \log(\mathcal{Z}(T, V, N)) \quad (2.0.2)$$

$$\mathcal{U}(T, V, N) = \mathcal{F} + TS = \mathcal{F} - T \frac{\partial \mathcal{F}}{\partial T} = k_B T^2 \frac{1}{\mathcal{Z}} \frac{\partial \mathcal{Z}}{\partial T} \quad (2.0.3)$$

where  $x_i \in \mathcal{V} \subset \mathbb{R}^3$  and  $p_i \in \mathbb{R}^3$ . The relation  $V = \text{vol}(\mathcal{V})$  describes the total volume occupied by the medium while  $T$  and  $N$  are its temperature and particle number respectively. To fully calculate  $\mathcal{U}$  it is necessary to obtain the partition function  $\mathcal{Z}$ . From there  $\mathcal{U}/V$  as the energy density can be compared to the pressure defined by

$$p = \frac{d\mathcal{F}}{dV} \quad (2.0.4)$$

to yield an EoS.

### 2.1 The Ultra-Relativistic Internal Energy

The ultra-relativistic case is a textbook example and well known limit that we will use in order to verify our later calculated results. First we write down the ultra-relativistic Hamiltonian given by

$$H(x, p) = ||p||c. \quad (2.1.1)$$

The corresponding partition function for a  $N$  particle system then reads ( $\beta = (k_B T)^{-1}$ )

$$\mathcal{Z} = \frac{V^N}{N! h^{3N}} \left[ \int_{\mathbb{R}^3} \exp(-\beta H(p)) d^3 p \right]^N \quad (2.1.2)$$

$$= \frac{V^N}{N! h^{3N}} \left[ \int_0^\infty 4\pi p^2 \exp(-\beta pc) dp \right]^N \quad (2.1.3)$$

$$= \frac{V^N}{N! h^{3N}} \frac{(4\pi)^N}{(\beta c)^{3N}} \left[ \int_0^\infty x^2 \exp(-x) dx \right]^N \quad (2.1.4)$$

$$\mathcal{Z} = \frac{1}{N!} \left( 8\pi V \left( \frac{k_B T}{hc} \right)^3 \right)^N \quad (2.1.5)$$

where from the first to second line we used usual spherical coordinates and afterwards the integral transformation  $x = \beta cp$ . The integral in equation (2.1.4) can be solved exactly with value 2. In the last line  $\beta = (k_B T)^{-1}$  was used for visual clarity. With equation (2.0.3) the internal energy  $\mathcal{U} = 3Nk_B T$  and with equation (2.0.4) the EoS can now be written down

$$p = \frac{Nk_B T}{V} = \frac{1}{3} \frac{\mathcal{U}}{V} = \frac{1}{3} \rho. \quad (2.1.6)$$

## 2.2 The Special-Relativistic Internal Energy

A source for the following calculations could not be found and were thus carried out by the author solely. The special relativistic Hamiltonian is given by

$$H = mc^2 \sqrt{1 + \frac{p^2}{m^2 c^2}}. \quad (2.2.1)$$

The ultra-relativistic limit can be obtained by letting  $m \rightarrow 0$ . In this limiting case we should be able to recover the results from equation (2.1.5).

$$\mathcal{Z} = \frac{V^N}{N! h^{3N}} \left[ \int_{\mathbb{R}^3} \exp \left( -\frac{mc^2 \sqrt{1 + \frac{p^2}{m^2 c^2}}}{k_B T} \right) d^3 p \right]^N \quad (2.2.2)$$

$$= \frac{V^N}{N! h^{3N}} \left[ \int_0^\infty 4\pi p^2 \exp \left( -\beta mc^2 \sqrt{1 + \frac{p^2}{m^2 c^2}} \right) dp \right]^N \quad (2.2.3)$$

$$= \frac{(4\pi V)^N}{N!} \left( \frac{mc}{h} \right)^{3N} \left[ \int_0^\infty q^2 \exp \left( -\alpha \sqrt{1 + q^2} \right) dq \right]^N \quad (2.2.4)$$

$$= \frac{(4\pi V)^N}{N!} \left( \frac{mc}{h} \right)^{3N} \left( \int_0^\infty \sinh(x)^2 \cosh(x) \exp(-\alpha \cosh(x)) dx \right)^N \quad (2.2.5)$$

$$= \frac{(4\pi V)^N}{N!} \left( \frac{mc}{h} \right)^{3N} \left( \int_0^\infty \frac{\sinh(2x)}{2} \sinh(x) \exp(-\alpha \cosh(x)) dx \right)^N \quad (2.2.6)$$

$$= \frac{(4\pi V)^N}{N!} \left( \frac{mc}{h} \right)^{3N} \left( -\frac{\sinh(2x)}{2\alpha} \exp(-\alpha \cosh(x)) \Big|_0^\infty \right. \quad (2.2.7)$$

$$\left. + \frac{1}{\alpha} \int_0^\infty \cosh(2x) \exp(-\alpha \cosh(x)) dx \right)^N \quad (2.2.8)$$

$$= \frac{(4\pi V)^N}{N!} \left( \frac{mc}{h} \right)^{3N} \left( \frac{1}{\alpha} \int_0^\infty \cosh(2x) \exp(-\alpha \cosh(x)) dx \right)^N \quad (2.2.9)$$

$$= \frac{1}{N!} \left( 8\pi V \left( \frac{k_B T}{hc} \right)^3 \frac{\alpha^2 K_2(\alpha)}{2} \right)^N \quad (2.2.10)$$

In the first step we used spherical coordinates followed by the substitution  $qmc = p$  and  $\alpha = \beta mc^2 = mc^2/k_B T$ . Afterwards we substituted  $q = \sinh(x)$  and used the identity  $\cosh(x) \sinh(x) = \sinh(2x)/2$ . Partial integration then leads to the last integral which can be identified as the modified Bessel function of the 2nd kind  $K_2(\alpha)$  [AS84]. The equation is then rewritten such that the ultra-relativistic limit can be read off upon letting  $\alpha \rightarrow 0$ . We can now calculate the internal energy  $\mathcal{U}$  from  $\mathcal{Z}$  via equation (2.0.3)

$$\mathcal{U} = 3Nk_B T - Nk_B T \alpha \left( \frac{\partial_\alpha K_2(\alpha)}{K_2(\alpha)} + 2 \right) \quad (2.2.11)$$

$$\mathcal{U} = 3Nk_B T - mc^2 \left( \frac{\partial_\alpha K_2(\alpha)}{K_2(\alpha)} + 2 \right). \quad (2.2.12)$$

Again, it can be seen that the ultra-relativistic limit can be obtained by letting  $\alpha \rightarrow 0$ , since the term written in parenthesis vanishes. It becomes immediately clear that the ultra-relativistic EoS is  $\rho = 3p$  which is equivalent to a traceless energy momentum tensor.

From equation (2.2.10) and (2.0.2), we immediately derive the ideal gas equation via the definition of pressure (see equation (2.0.4) in the canonical ensemble

$$p = \frac{\partial \mathcal{F}}{\partial V} = \frac{Nk_B T}{V}. \quad (2.2.13)$$

Note that in the case of the Hamiltonian  $H$  having a non-trivial dependence on the coordinates  $x_i$ , the partition function would differ by a factor involving  $V$  and  $T$ . This more general scenario does not yield the ideal gas law and would thus need to be treated separately.

## 2.3 The Special-Relativistic Equation of State

This section aims to develop an equation between the thermodynamic energy density  $\rho = \mathcal{U}/V$  and the pressure  $p$  of the gas given by the ideal gas equation (2.0.4). We assume an additional constraint namely adiabaticity and thus further reduce the degrees of freedom of the thermodynamic system.

When assuming an adiabatic condition  $\delta Q = 0$  and using the First Law of Thermodynamics [Fli18]  $dU = \delta Q + \delta W$ , where  $\delta W = -p dV$  and  $dU = C_V dT$ , we can relate pressure and temperature. This adiabatic condition is quite well satisfied since the timescale of radiational and other losses compared to the timescale of thermodynamic events is negligible [Noe08; Vin17]. Using equation (2.2.12) and (2.0.4), we obtain

$$-p dV = C_V dT \quad (2.3.1)$$

$$-\frac{Nk_B T}{V} dV = Nk_B \left[ 1 + \alpha^2 \left( \left( \frac{\partial_\alpha K_2(\alpha)}{K_2(\alpha)} \right)^2 - \frac{\partial_\alpha^2 K_2(\alpha)}{K_2(\alpha)} \right) \right] dT \quad (2.3.2)$$

$$-\frac{dV}{V} = \left[ 1 - \alpha^2 \partial_\alpha \left( \frac{\partial_\alpha K_2(\alpha)}{K_2(\alpha)} \right) \right] \frac{dT}{T} \quad (2.3.3)$$

$$= \left[ 1 - \alpha^2 \partial_\alpha^2 (\log K_2(\alpha)) \right] \frac{dT}{T} \quad (2.3.4)$$

This equation also shows explicitly the  $T$  dependence of the specific heat  $C_V$  for the general case. Again, taking the ultra-relativistic limit by letting  $\alpha \rightarrow 0$ , the right-hand

term in the first equation converges to  $-2Nk_B$ . This agrees with the expected specific heat for an ultra-relativistic gas  $C_{V,\text{ur}} = 3Nk_B$ . With the identity  $d\alpha/\alpha = -dT/T$  (using  $\alpha = mc^2/k_B T$ ), we can transform the equation and integrate it. After applying partial integration, the result is

$$\frac{dV}{V} = (1 - \alpha^2 \partial_\alpha^2 \log K_2(\alpha)) \frac{d\alpha}{\alpha} \quad (2.3.5)$$

$$\log\left(\frac{V}{V_0}\right) = \log\left(\frac{\alpha}{\alpha_0}\right) - \int_{\alpha_0}^{\alpha} \alpha \frac{\partial^2}{\partial \alpha^2} \log(K_2(\alpha')) d\alpha' \quad (2.3.6)$$

$$= \log\left(\frac{\alpha}{\alpha_0}\right) + \log\left(\frac{K_2(\alpha)}{K_2(\alpha_0)}\right) - \left[\alpha \frac{\partial_\alpha K_2}{K_2}\right]_{\alpha_0}^{\alpha} \quad (2.3.7)$$

This equation enables us to write down a relation between volume and temperature (encapsulated in  $\alpha = mc^2/k_B T$ )

$$V(\alpha) = \frac{\alpha K_2(\alpha)}{C} \exp\left(\alpha \frac{K_3(\alpha) + K_1(\alpha)}{2K_2(\alpha)}\right) \quad (2.3.8)$$

where the constant  $C$  is defined by the equation beforehand and only depends on the integration boundaries  $\alpha_0$  and  $V_0$ . It is given by

$$C = \frac{\alpha_0 K_2(\alpha_0)}{V_0} \exp\left(\alpha_0 \frac{K_1(\alpha_0) + K_3(\alpha_0)}{2K_2(\alpha_0)}\right). \quad (2.3.9)$$

Since the goal of this section is to obtain a readable output for an EoS, it is necessary to construct a bijection relating  $p$  and  $T$ . This becomes clear when writing down the energy density

$$\rho = \frac{\mathcal{U}}{V} = \frac{Nk_B T}{V} - \frac{Nk_B T}{V} \left(\alpha \frac{\partial_\alpha K_2(\alpha)}{K_2(\alpha)}\right) \quad (2.3.10)$$

where  $p = Nk_B T/V$  can be easily identified but the  $T$  dependence via  $\alpha$  is not solved yet. The pressure  $p$  can be rewritten to take the form

$$p = \frac{Nk_B T}{V} = C N m c^2 \frac{1}{K_2(\alpha) \alpha^2} \exp\left(-\alpha \frac{K_1(\alpha) + K_3(\alpha)}{2K_2(\alpha)}\right). \quad (2.3.11)$$

At this point it is not reasonable to ask what happens in the ultra-relativistic limit since  $C$  depends non-trivially on  $m$  and thus  $m$  is not fully substituted in  $\alpha$ .

Interestingly, the pressure seems to be constant for very high temperatures. The limiting case is obtained when taking  $T \rightarrow \infty$  (which corresponds to  $\alpha \rightarrow 0$ )

$$\lim_{\alpha \rightarrow 0} \left[ \frac{1}{K_2(\alpha) \alpha^2} \exp\left(-\alpha \frac{K_1(\alpha) + K_3(\alpha)}{2K_2(\alpha)}\right) \right] = \frac{1}{2e^2} \approx 0.006767. \quad (2.3.12)$$

The same argument then holds true for the density given by equation (2.3.10) and since

$$\lim_{\alpha \rightarrow 0} \left[ 1 + \alpha \frac{K_1(\alpha) + K_3(\alpha)}{2K_2(\alpha)} \right] = 3 \quad (2.3.13)$$

we have

$$\lim_{\alpha \rightarrow 0} \left( \frac{\rho(\alpha)}{C m c^2} \right) = \frac{3}{2e^2} \approx 0.203003. \quad (2.3.14)$$



### Theorem 2.1

The mapping  $p : \mathbb{R}_{>0} \rightarrow \mathbb{R}_{>0}, \alpha \mapsto p(\alpha)$  written down in equation (2.3.11) is a bijection for any  $N, m, c, C \neq 0$ .

$$p = \frac{Nk_B T}{V} = CNmc^2 \frac{1}{K_2(\alpha)\alpha^2} \exp\left(-\alpha \frac{K_1(\alpha) + K_3(\alpha)}{2K_2(\alpha)}\right) \quad (2.3.15)$$

*Proof.* For this proof it suffices to show that the function  $p(\alpha)$  has a strictly monotonous behaviour. Without loss of generality, we assume  $N, m, c, C > 0$ . Now it is obvious that the first two terms  $Nmc^2/C K_2(\alpha)$  and  $\alpha^{-2}$  are strictly decreasing. This is easy to see when using [AS84]

$$\frac{\partial K_n}{\partial \alpha} = \frac{n}{\alpha} K_n - K_{n+1} = -\frac{K_{n-1} + K_{n+1}}{2}. \quad (2.3.16)$$

We then calculate the derivative of the third term and divide by the exponential (since it is positive)

$$\frac{1}{\exp(\alpha \partial_\alpha \log(K_2))} \frac{\partial}{\partial \alpha} \exp(\alpha \partial_\alpha \log(K_2)) \quad (2.3.17)$$

$$= \partial_\alpha \log(K_2) + \alpha \partial_\alpha^2 \log(K_2) \quad (2.3.18)$$

$$= \frac{\partial_\alpha K_2}{K_2} + \alpha \frac{\partial_\alpha^2 K_2}{K_2} + \alpha \left( \frac{\partial_\alpha K_2}{K_2} \right)^2 \quad (2.3.19)$$

$$= \frac{K_1 - \frac{2}{\alpha} K_2}{K_2} + \alpha \frac{\frac{1}{\alpha} K_1 - K_2 - \frac{2}{\alpha} (K_1 - \frac{2}{\alpha} K_2) + \frac{2}{\alpha^2} K_2}{K_2^2} + \alpha \frac{K_1^2 - \frac{1}{\alpha} K_1 K_2 - \frac{4}{\alpha^2} K_2^2}{K_2^2} \quad (2.3.20)$$

$$= -\alpha + \alpha \frac{K_1^2}{K_2^2} - 4 \frac{K_1}{K_2} \quad (2.3.21)$$

thus, it is sufficient to show that

$$\alpha \frac{K_1^2}{K_2^2} < \alpha + 4 \frac{K_1}{K_2}. \quad (2.3.22)$$

We quickly prove the more general result  $K_\nu < K_{\nu+1}$ . One possible definition [AS84] for the Bessel function  $K_\nu$  is given by

$$K_\nu := \frac{\sqrt{\pi}}{(\nu - \frac{1}{2})!} \left( \frac{1}{2} z \right)^\nu \int_1^\infty e^{-tz} (t^2 - 1)^{\nu - \frac{1}{2}} dt \quad (2.3.23)$$

We inspect the ratio

$$\frac{K_\nu}{K_{\nu+1}} = \frac{1}{(\nu + \frac{1}{2}) (\frac{1}{2} z)} \frac{\int e^{-tz} (t^2 - 1)^{\nu - 1/2} dt}{\int e^{-tz} (t^2 - 1)^{\nu + 1/2} dt} \quad (2.3.24)$$

and rewrite the denominator with partial integration

$$\frac{1}{2} z \int_1^\infty e^{-tz} (t^2 - 1)^{\nu + 1/2} dt = \left( \nu + \frac{1}{2} \right) \int_1^\infty e^{-tz} t (t^2 - 1)^{\nu - 1/2} dt. \quad (2.3.25)$$

Now it is obvious that  $K_{\nu+1} > K_\nu$ . Thus in total, the function given by equation (2.3.26) can be inverted.  $\square$

With this mapping  $p : \mathbb{R}_{>0} \rightarrow \mathbb{R}_{>0}$ ,  $\alpha \mapsto p(\alpha)$  and its inverse  $\alpha : \mathbb{R}_{>0} \rightarrow \mathbb{R}_{>0}$ ,  $p \mapsto \alpha(p)$ , we can use (2.3.10) and finally write down the EoS

$$\rho = \frac{\mathcal{U}}{V} = p \left( 1 + \alpha(p) \frac{K_1(\alpha(p)) + K_3(\alpha(p))}{2K_2(\alpha(p))} \right). \quad (2.3.26)$$

Figure 1 is obtained by choosing numerical values and then interpolating and inverting equation (2.3.11). The constant factor

$$B := CNmc^2/p_0 \quad (2.3.27)$$

is substituted to obtain independence of  $p_0$ <sup>1</sup>. Furthermore the graphs of the plotted EoS are normalised such that

$$\rho_i(p_0) = \rho_{0,i} \quad (2.3.28)$$

and thus can be compared with each other. Note that the  $\rho_0$  of the plot is not a universal value but rather each function has been scaled individually and  $\rho_0$  is a placeholder for the corresponding  $\rho_{0,i}$ . The polytropic EoS  $\rho_{cla}(p) = Ap^{1/\gamma}$  when normalised as before, the equation is independent of  $A$  and can be uniquely characterised by  $n$ . For further details, the interested reader is referred to [Ple21].

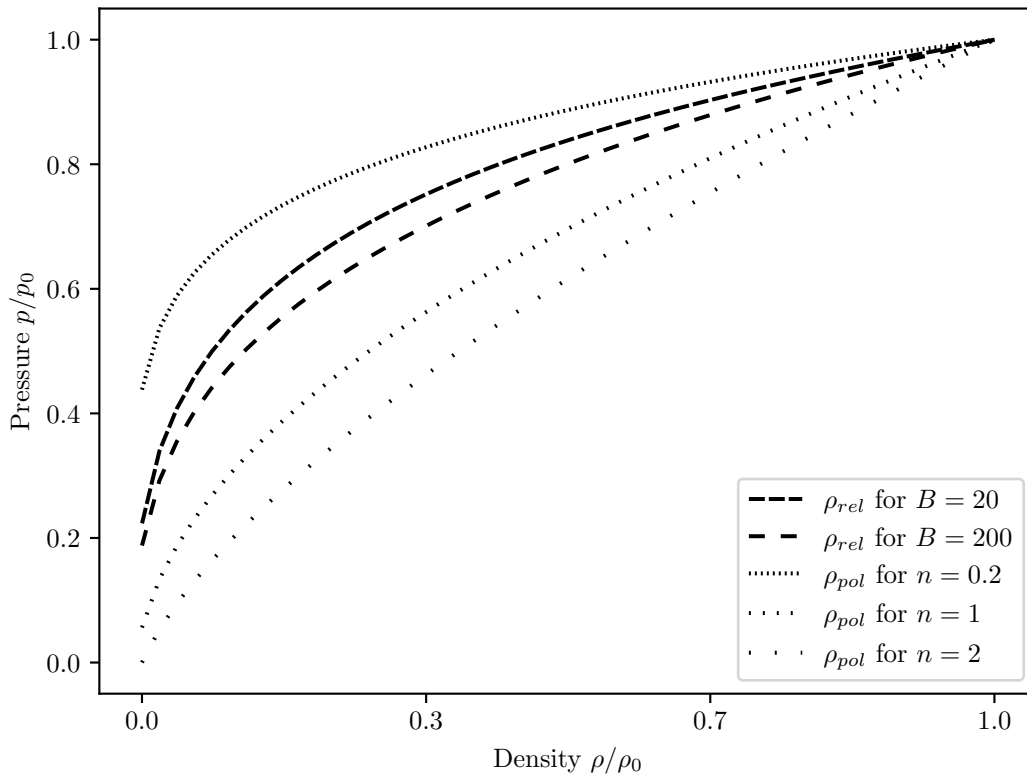


Figure 1: Relativistic Equation of State

The relativistic EoS  $\rho(p)$  is normalised such that values can be compared with a polytropic EoS. Graphs for the relativistic version are independent of the exponent  $n$  which is a degree of freedom intrinsic to the polytropic EoS. By normalisation, the graphs of the polytropic EoS are independent of the factor  $A$ .

<sup>1</sup>This will be the starting value for the pressure in the Tollmann-Oppenheimer-Volkoff equation. Since the pressure decreases from the inside a star to the outside, this resembles our highest value of  $p$ .

### 3 Calculating the Mass of a Star with an EoS

#### 3.1 Deriving the TOV-Equation

This chapter derives the TOV equation which was introduced earlier in Information about GR can be found in [CD00; Cho09; Cho15; Wal84]. We follow the derivation of Wald. We consider a spherical-symmetric static Lorentz-Manifold  $(V, g)$  with charts such that the metric  $g$  can be written as

$$g = -e^\nu dt^2 + e^\lambda dr^2 + r^2(d\theta^2 + \sin^2 \theta d\phi^2). \quad (3.1.1)$$

The stress-energy tensor of an ideal fluid with density  $\rho$  and pressure  $p$  is given by

$$T_{\mu\nu} = \rho u_\mu u_\nu + p(g_{\mu\nu} + u_\mu u_\nu) \quad (3.1.2)$$

where  $u$  is the 4-velocity of the fluid. In the rest frame where  $u^\mu = (-e^{-\nu/2}, 0, 0, 0)$ , this equation simplifies to

$$T_\nu^\mu = \text{diag}(-\rho, p, p, p) \quad (3.1.3)$$

The Christoffel symbols for this metric are

$$\Gamma_{\mu\nu}^0 = \begin{bmatrix} 0 & \nu'/2 & & \\ \nu'/2 & 0 & & \\ & & 0 & \\ & & & 0 \end{bmatrix}, \quad \Gamma_{\mu\nu}^1 = \begin{bmatrix} \nu' e^{\nu-\lambda}/2 & & & \\ & \lambda'/2 & & \\ & & -r e^{-\lambda} & \\ & & & -r \sin^2 \theta e^{-\lambda} \end{bmatrix} \quad (3.1.4)$$

$$\Gamma_{\mu\nu}^2 = \begin{bmatrix} 0 & & & \\ & 0 & 1/r & \\ & 1/r & 0 & \\ & & & -\sin \theta \cos \theta \end{bmatrix}, \quad \Gamma_{\mu\nu}^3 = \begin{bmatrix} 0 & 0 & & \\ 0 & 0 & 1/r & \\ 0 & 0 & \cos \theta / \sin \theta & \\ 1/r & \cos \theta / \sin \theta & 0 & \end{bmatrix} \quad (3.1.5)$$

From these, the non-zero components of the Ricci-Tensor can be calculated

$$R_{11} = \frac{1}{4r} e^{-\lambda} [(2r\nu'' + r\nu'^2) + (4 - r\lambda')\nu'] \quad (3.1.6)$$

$$R_{22} = -\frac{1}{4r} e^{-\lambda} [(2r\nu'') + r\nu'^2 - r\lambda'\nu' - 4\lambda'] \quad (3.1.7)$$

$$R_{33} = -\frac{1}{2r^2} e^{-\lambda} (r\nu' - r\lambda' - 2e^\lambda + 2) \quad (3.1.8)$$

$$R_{44} = R_{33} \quad (3.1.9)$$

and with  $R_{\mu\nu} - g_{\mu\nu}R/2 = G_{\mu\nu} = 8\pi T_{\mu\nu}$  we ultimately receive the following field equations.

$$-8\pi T_0^0 = 8\pi\rho = \frac{\lambda' e^{-\lambda}}{r} + \frac{1 - e^{-\lambda}}{r^2} \quad (3.1.10)$$

$$8\pi T_1^1 = 8\pi p = \nu' \frac{e^{-\lambda}}{r} - \frac{1 - e^{-\lambda}}{r^2} \quad (3.1.11)$$

$$8\pi T_2^2 = 8\pi p = \frac{e^{-\lambda}}{2} \left[ \nu'' + \left( \frac{\nu'}{2} + \frac{1}{r} \right) (\nu' - \lambda') \right] \quad (3.1.12)$$

Since  $R_4^4 = R_3^3$ , we omitted the last equation. From equation (3.1.10) we infer the relation.

$$e^{-\lambda} = 1 - \frac{2}{r} \int_0^r 4\pi\rho(r')r'^2 dr' = 1 - \frac{2m(r)}{r}. \quad (3.1.13)$$

The metric needs to be defined at every point in space and thus can not have any additional integration constant in equation (3.1.13), since otherwise we would obtain a term  $a/r$  which is not defined for  $r \rightarrow 0$ .

The property  $m(r)$  can be recognised as the Newtonian mass of the star (which is different to the proper mass [Wal84]). Since  $e^{-\lambda} > 0$ , we immediately see that  $m(r) < r/2$ .

In addition to the Field equations (3.1.10) through (3.1.12) the divergence of the Stress-Energy Tensor yields more information

$$\nabla_\mu T^{\mu\nu} = 0. \quad (3.1.14)$$

The following explicit calculation, which again assumes spherical symmetry, shows how to obtain this additional restriction on the pressure and density.

$$\nabla_\mu T^\mu_\nu = \partial_\mu T^\mu_\nu + \Gamma^\mu_{\mu\sigma} T^\sigma_\nu - \Gamma^\sigma_{\mu\nu} T^\mu_\sigma \quad (3.1.15)$$

$$\nabla_\mu T^\mu_1 = \frac{\partial p}{\partial r} + p \left( \Gamma^0_{01} + \Gamma^1_{11} + \Gamma^2_{21} + \Gamma^3_{31} \right) - \Gamma^\sigma_{\mu 1} T^\mu_\sigma \quad (3.1.16)$$

$$= \frac{\partial p}{\partial r} + p \left( \frac{\nu' + \lambda'}{2} + \frac{2}{r} \right) + \rho \frac{\nu'}{2} - p \frac{\lambda'}{2} - p \frac{2}{r} \quad (3.1.17)$$

$$\frac{\partial p}{\partial r} = -\frac{p + \rho}{2} \nu' \quad (3.1.18)$$

Together with equation (3.1.11) and the definition (3.1.13), we can write

$$\frac{\partial p}{\partial r} = -\frac{p + \rho}{2} \left( \frac{8\pi p r + \frac{1-e^{-\lambda}}{r}}{e^{-\lambda}} \right) \quad (3.1.19)$$

$$= -\frac{p + \rho}{2r} \left( \frac{8\pi p r + \frac{2m}{r^2}}{1 - \frac{2m}{r}} \right) \quad (3.1.20)$$

$$= -\frac{m\rho}{r^2} \left( 1 + \frac{p}{\rho} \right) \left( \frac{4\pi r^3 p}{m} + 1 \right) \left( 1 - \frac{2m}{r} \right)^{-1} \quad (3.1.21)$$

$$\frac{\partial p}{\partial r} = -\frac{Gm\rho}{r^2} \left( 1 + \frac{p}{\rho c^2} \right) \left( \frac{4\pi r^3 p}{mc^2} + 1 \right) \left( 1 - \frac{2Gm}{rc^2} \right)^{-1} \quad (3.1.22)$$

where in the last step the constants  $c = G = 1$  were put back in. Equation (3.1.22) together with (3.1.13) yields the TOV differential equations

$$\frac{\partial m}{\partial r} = 4\pi \rho(r) r^2 \quad (3.1.23)$$

$$\frac{\partial p}{\partial r} = -\frac{Gm\rho}{r^2} \left( 1 + \frac{p}{\rho c^2} \right) \left( \frac{4\pi r^3 p}{mc^2} + 1 \right) \left( 1 - \frac{2Gm}{rc^2} \right)^{-1} \quad (3.1.24)$$

### 3.2 Newtonian Limit

This section follows [Wei20] and [Cha58, pp. 89 sqq.]. Together with a polytropic EoS  $p = K\rho^{1+1/n}$  and the definition  $\rho = \lambda\theta^n$ , we expect to obtain the Newtonian behavior in the non-relativistic limit in the form of the Lane-Emden (LE) equation

$$\frac{K(n+1)\lambda^{1/n-1}}{4\pi}\Delta\theta + \theta^n = 0. \quad (3.2.1)$$

The usual non-relativistic limit is obtained from a Taylor expansion of equation (3.1.24) around  $1/c^2$  in lowest order. The resulting equation then reads

$$\frac{\partial p}{\partial r} = -\frac{Gm\rho}{r^2} + \mathcal{O}\left(\frac{1}{c^2}\right). \quad (3.2.2)$$

One could be tempted to make the assumption that  $\frac{\partial p_{\text{TOV}}}{dr} \leq \frac{\partial p_{\text{LE}}}{dr}$ . However this equation fails since the mass given in equation (3.1.24) is not the same as the one given in the LE equation. Using the previous relations for  $\rho$  and  $p$  and again setting  $G = c = 1$ , we can calculate

$$\frac{\partial p}{\partial r} = \frac{\partial}{\partial r}(K\rho^{1+1/n}) = K\lambda^{1+1/n}(n+1)\theta^n \frac{\partial \theta}{\partial r} = -\frac{m\lambda\theta^n}{r^2} \quad (3.2.3)$$

by using the definition of our polytropic EoS and equation (3.2.2). Rearranging and taking the derivative of this equation and using  $\partial m/\partial r = 4\pi\rho r^2$ , we obtain

$$-\frac{\partial m}{\partial r} = K\lambda^{1/n}(n+1)\frac{\partial}{\partial r}\left(r^2\frac{\partial \theta}{\partial r}\right) = -4\pi r^2\lambda\theta^n \quad (3.2.4)$$

Upon redefining  $\xi = r/\kappa$  where  $4\pi\kappa^2 = (n+1)K\lambda^{1/n-1}$ , one can obtain the mathematically cleaner looking equation

$$\frac{1}{\xi^2}\frac{\partial}{\partial \xi}\left(\xi^2\frac{\partial \theta}{\partial \xi}\right) + \theta^n = 0 \quad (3.2.5)$$

Exact solutions are known for the cases  $n = 0, 1, 5$ . The derivation can be found in appendix A. Figure 2 and Table 1 summarise them. Together with equation (3.2.5), we can already suspect that for values  $n \geq 5$ , the equation does not yield solutions with a zero value.

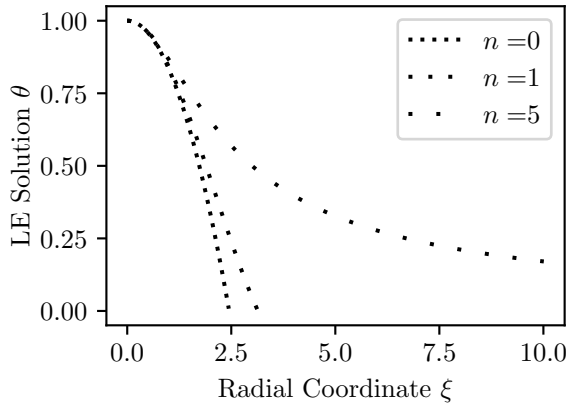


Figure 2: Graph of exact LE solutions.

$n$	LE Solution	$\xi_0$
$n = 0$	$1 - \frac{1}{6}\xi^2$	$\sqrt{6}$
$n = 1$	$\frac{\sin(\xi)}{\xi}$	$\pi$
$n = 5$	$\frac{1}{\sqrt{1 + \frac{1}{3}\xi^2}}$	$\infty$

Table 1: LE exact Solutions and their zero value.

### 3.3 Mass Bounds

We will first follow the approach given in [Wal84]. The first assumptions will be  $d\rho/dr < 0$  and  $\rho \geq 0$ . Also we again consider a compact star, meaning  $\rho(r) = 0$  for all  $r > R$ . We first state a useful Lemma for the proof of our next theorem.

**Lemma 3.1**

Given  $\rho \in C^1(\mathbb{R}_{\geq 0})$  is monotonously decreasing, the function  $\rho_{av} = m/r^3$  has negative slope.

*Proof.* Taking the derivative of  $\rho_{av}$  and applying equation (3.1.13), we obtain

$$\frac{\partial}{\partial r} \left( \frac{m}{r^3} \right) = -3 \frac{m}{r^4} + \frac{4\pi\rho}{r}. \quad (3.3.1)$$

By equation (3.1.24) we know that  $p$  has negative slope and since  $\rho$  is monotonously increasing, we have

$$m = \int_0^r 4\pi\rho r^2 dr \geq \rho \int_0^r 4\pi r^2 = \frac{4\pi r^3 \rho}{3} \quad (3.3.2)$$

which shows that  $\frac{4\pi\rho r^3}{3} \leq m$  and completes the proof.  $\square$

**Theorem 3.2 - Mass Bound**

The Mass of a spherically symmetric star is bound from above by

$$M < \frac{4}{9}R. \quad (3.3.3)$$

*Proof.* In our attempt to obtain an upper limit for the mass of a spherically symmetric star, we start by taking the difference of equation (3.1.11) and (3.1.12), we obtain

$$0 = \nu' \frac{e^{-\lambda}}{r} - \frac{1 - e^{-\lambda}}{r^2} - \frac{e^{-\lambda}}{2} \left[ \nu'' + \left( \frac{\nu'}{2} + \frac{1}{r} \right) (\nu' - \lambda') \right] \quad (3.3.4)$$

$$= -\frac{2m(r)}{r^3} + \frac{\lambda' e^{-\lambda}}{2r} - \frac{e^{-\lambda}}{2} \left[ \nu'' + \frac{\nu'^2}{2} - \frac{\nu'}{r} - \frac{\lambda' \nu'}{2} \right] \quad (3.3.5)$$

$$= r \frac{\partial}{\partial r} \left( \frac{m(r)}{r^3} \right) - \frac{e^{-\lambda}}{2} \left[ \nu'' + \frac{\nu'^2}{2} - \frac{\nu'}{r} - \frac{\lambda' \nu'}{2} \right] \quad (3.3.6)$$

$$0 = \frac{\partial}{\partial r} \left( \frac{m(r)}{r^3} \right) - \frac{e^{-\lambda}}{2} \left[ \frac{\nu''}{r} + \frac{\nu'^2}{2r} - \frac{\nu'}{r^2} - \frac{\lambda' \nu'}{2r} \right] \quad (3.3.7)$$

$$= \frac{\partial}{\partial r} \left( \frac{m(r)}{r^3} \right) - \frac{1}{2} e^{-\frac{\lambda+\nu}{2}} \frac{\partial}{\partial r} \left[ \frac{1}{r} \nu' e^{\frac{\nu-\lambda}{2}} \right]. \quad (3.3.8)$$

Since  $\partial_r \rho \leq 0$ , also the average density  $m(r)/r^3$  decreases with  $r$ . Thus we obtain

$$\frac{\partial}{\partial r} \left[ \frac{1}{r} \nu' \exp \left( \frac{\nu - \lambda}{2} \right) \right] \leq 0. \quad (3.3.9)$$

We integrate this equation from  $R$  to  $r < R$

$$\frac{\nu'}{r} \exp \left( \frac{\nu - \lambda}{2} \right) \geq \frac{2\nu'(R)}{R} e^{-\frac{1}{2}\lambda(R)} \frac{\partial}{\partial r} e^{\frac{\nu}{2}} \Big|_R \quad (3.3.10)$$

and use the Schwarzschild solution at  $r = R$  for  $e^\lambda$  and  $e^\nu$ . This is justified since we assumed  $\rho(r) = 0$  for  $r > R$  and thus we need to recover the vacuum solution for a spherically symmetric object which is given by the Schwarzschild solution. By continuity of the metric on every point of space, we can match

$$e^{-\lambda(r)}|_R = \left[1 - \frac{2M}{r}\right]_R = e^{\nu(r)}|_R \quad (3.3.11)$$

and with the explicit solution for  $e^{-\lambda}$ , we obtain

$$\frac{2m(r)}{r}\bigg|_R = \frac{2M}{R}. \quad (3.3.12)$$

When plugging this into equation (3.3.11), the result is

$$\frac{\nu'}{2r} \exp\left(\frac{\nu - \lambda}{2}\right) \geq \frac{(1 - 2M/R)^{1/2}}{R} \frac{\partial}{\partial r} \left(1 - \frac{2M}{r}\right)^{1/2} \bigg|_{r=R} = \frac{M}{R^3}. \quad (3.3.13)$$

Now we multiply by  $r \exp(\lambda/2)$  and use the explicit solution for  $e^\lambda$

$$\frac{\partial}{\partial r} (e^{\frac{\nu}{2}}) \geq \frac{M}{R^3} r e^{\frac{\lambda}{2}} = \frac{M}{R^3} (r - 2m(r)) \quad (3.3.14)$$

and integrate again this time from 0 to  $R$

$$e^{\nu(0)/2} \leq \left(1 - \frac{2M}{R}\right)^{1/2} - \frac{M}{R^3} \int_0^R \left[1 - \frac{2m(r)}{r}\right]^{-1/2} r dr. \quad (3.3.15)$$

As we have already noted, the average density  $m(r)/r^3$  decreases, meaning explicitly  $m(r)/r^3 \geq M/R^3$  and thus the integral with the previous equation can be written as

$$e^{\nu(0)/2} \leq \left(1 - \frac{2M}{R}\right)^{1/2} + \frac{1}{2} \left[1 - \frac{2Mr^2}{R^3}\right]^{1/2} \bigg|_{r=0}^{r=R} = \frac{3}{2} \left(1 - \frac{2M}{R}\right)^{1/2} - \frac{1}{2}. \quad (3.3.16)$$

The simple fact that  $e^{\nu(0)/2} > 0$  then implies

$$\left(1 - \frac{2M}{R}\right)^{1/2} > \frac{1}{3} \quad (3.3.17)$$

which is equivalent to

$$M < \frac{4R}{9}. \quad (3.3.18)$$

□

This shows that the mass of star has an upper limit under the assumptions  $\rho \geq 0$ ,  $\partial_r \rho \leq 0$  and  $\rho(R) = 0$  for some  $R \geq 0$ .

## 4 Numerical Solutions

### 4.1 Comparing TOV and LE results with a polytropic EoS

In this section the numerical solutions of the TOV equation

$$\frac{\partial m}{\partial r} = 4\pi\rho r^2 \quad (4.1.1)$$

$$\frac{\partial p}{\partial r} = -\frac{m\rho}{r^2} \left(1 + \frac{p}{\rho}\right) \left(\frac{4\pi r^3 p}{m} + 1\right) \left(1 - \frac{2m}{r}\right)^{-1} \quad (4.1.2)$$

as derived previously in section 3.1 will be discussed. To obtain numerical solvability an EoS in the form  $\rho(r, p)$  is supplied. In Figure 3, a plot of such a solution is presented. The density  $\rho$  is derived via the equation (4.1.1) and the integration is done with a 4th order Runge-Kutta Method [Run95; Kut01; HSW10]. The integration is stopped once the pressure reaches values  $p \leq 0$ . The initial value of (4.1.1) is  $\partial_r m(r=0) = 0$ . For equation (4.1.1), the initial value can be calculated when applying L'Hôpital's rule and combining them to obtain

$$\lim_{r \rightarrow 0} \frac{m}{r} = \lim_{r \rightarrow 0} \frac{\partial m}{\partial r} = \lim_{r \rightarrow 0} \frac{4\pi\rho r^2}{1} = 0 \quad (4.1.3)$$

$$\lim_{r \rightarrow 0} \frac{m}{r^2} = \lim_{r \rightarrow 0} \frac{1}{2r} \frac{\partial m}{\partial r} = \lim_{r \rightarrow 0} \frac{4\pi\rho r^2}{2r} = 0 \quad (4.1.4)$$

$$\lim_{r \rightarrow 0} \frac{m}{r^3} = \lim_{r \rightarrow 0} \frac{1}{3r^2} \frac{\partial m}{\partial r} = \lim_{r \rightarrow 0} \frac{4\pi\rho r^2}{3r^2} = \frac{4\pi\rho_0}{3} \quad (4.1.5)$$

$$\lim_{r \rightarrow 0} \frac{\partial p}{\partial r} = 0 \quad (4.1.6)$$

Explicit code can be found in [Ple21]. The Lane-Emden equation was obtained in section 3.2 as the non-relativistic limit of the TOV equation by neglecting terms of order  $1/c^2$  and higher and setting  $G = c = 1$ . To obtain numerical results for the LE equation as given in equation (3.2.5), a substitution of the form  $d\theta/d\xi = \chi$  was used.

$$\begin{aligned} \frac{d\theta}{d\xi} &= \chi & \frac{d\theta}{d\xi} \Big|_{\xi=0} &= 0 \\ \frac{d\chi}{d\xi} &= -\frac{2}{\xi}\chi - \theta^n & \frac{d\chi}{d\xi} \Big|_{\xi=0} &= -\theta^n \end{aligned} \quad (4.1.7)$$

The initial value for  $d\chi/d\xi$  can be calculated with L'Hôpital's rule. With the conversion factor  $\kappa$  derived in section 3.2, TOV and LE results can be compared. Figure 3 shows the solution of both equations for the parameters of Table 2. Additionally conversion equations to compare TOV and LE results are displayed. Since the mass reads

$$m(r) = \int_0^r 4\pi r'^2 \rho(r') dr', \quad (4.1.8)$$

we expect  $\partial m/\partial r(R) = 0$  if  $p(R) = 0$  when choosing a polytropic equation of state with  $\gamma > 0$ . The plot in Figure 3 shows this expected behaviour for the Lane Emden equation at  $r \approx 2.31$  and has the same behaviour for the TOV results at  $r \approx 6.80$ .<sup>2</sup>

---

<sup>2</sup>For the purpose of nicely displaying the calculated result, the plot only shows result up to  $r = 2.5$



TOV		LE	
EOS	$\rho = A p^{1/\gamma}$	EOS	$p = K \rho^\gamma$
$A$	2		
$\gamma = 1 + \frac{1}{n}$	4/3	$n = 1/(\gamma - 1)$	3
$p_0$	0.5	$\theta_0$	1
$m_0$	0	$d\theta_0$	0
$dr$	0.01	$d\xi = dr/\kappa$	0.01/0.3355 $\approx$ 0.0298
$\rho_0 = A p_0^{1/\gamma}$	$2(2)^{\frac{4}{3}} \approx 1.1892$	$\lambda = \rho_0$	$2(2)^{\frac{4}{3}} \approx 1.1892$
		$K = A^{-1/\gamma}$	$2^{-3/4} \approx 0.5946$
		$\kappa^2 = ((n+1)K\lambda^{1/n-1})/(4\pi)$	$\approx 0.1125$

Table 2: Numerical Parameters for TOV and LE equation  
To compare results, values  $\rho_0, K, \kappa$  are calculated from ones supplied to the solving routine.

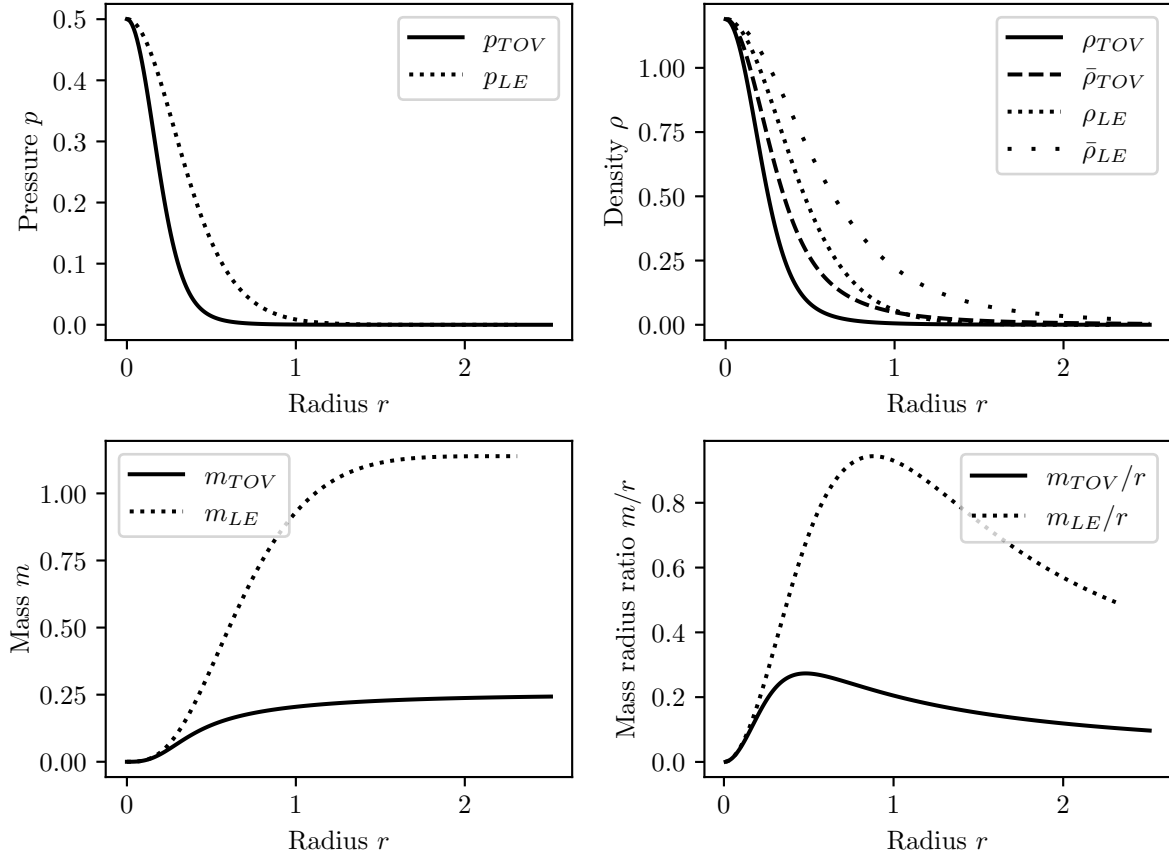


Figure 3: Comparison of the TOV and LE equation.  
The images show the plots for the parameters of Table 2. The pressures and masses as presented in the first column are direct solutions of the TOV or LE equations. On the top right-hand side the densities calculated with the EoS and average densities  $\bar{\rho}_i = (4\pi/3)^{-1}m_i/r^3$  are being compared.

## 4.2 Verifying the results

One can compare calculated LE results with already known exact solutions for certain exponents as given in table 1. The graphs of figure 4 overall show good numerical agreement. The integration stepsize for the left-hand side results of figure 4 is  $d\xi = 0.03$ . Initial values at  $r = 0$  are identical since they were set to be. The spike occurring afterwards can be explained by equations (4.1.7). The initial value of all derivatives is identically 0 which means no change in the values  $\theta$  or  $\chi$  for the initial step of the numerical integration. On the other hand the exact results will show a change which explains the large discrepancy in the beginning. Despite this behaviour we can see that the criterion for convergence is well fulfilled in the right-hand side of figure 4. The plot displays nicely that by lowering the integration stepsize, the maximum difference in this interval decreases.

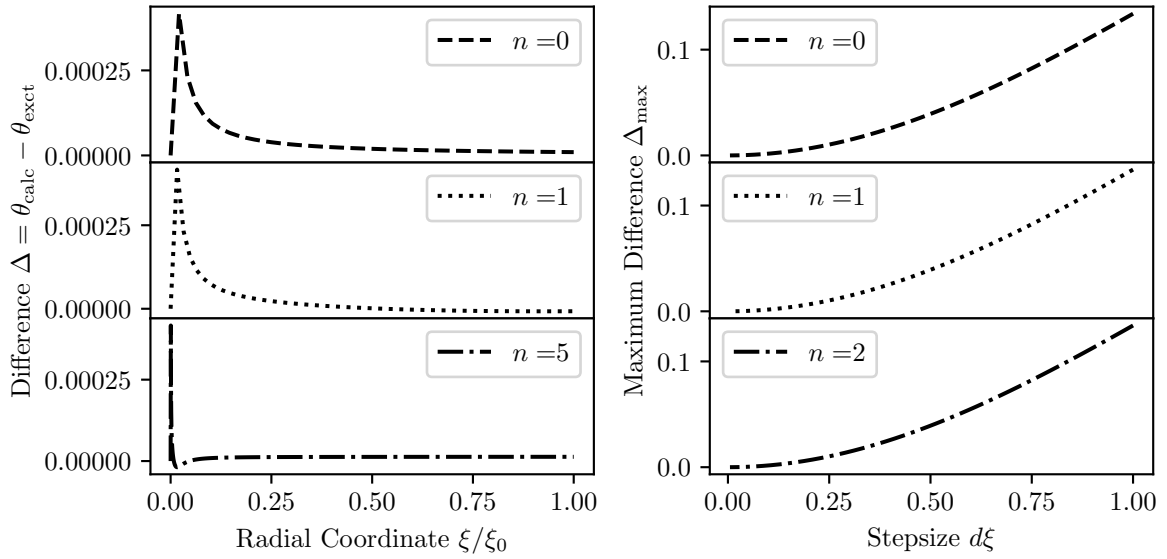


Figure 4: Validation of numerical LE results

The first two plots show absolute and relative difference of known exact and numerically calculated results while the third plot shows the behaviour of the maximum value of  $\Delta$  as the stepsize decreases.

## 4.3 Relativistic EoS

In the previous discussion, we relied on the EoS given in Table 2. This is a versatile assumption, but one could ask what would happen to a star in which the particles have no interaction but are close to relativistic speed. The resulting EoS was calculated in the beginning (see equation (2.3.26)) although not written down explicitly. Since explicit inversion of the given function is complicated, we rely on numerical methods for calculation. Section 2.3 showed that the function given in equation (2.3.11)

$$p(\alpha) = CNmc^2 \frac{1}{K_2(\alpha)\alpha^2} \exp\left(-\alpha \frac{K_1(\alpha) + K_3(\alpha)}{2K_2(\alpha)}\right). \quad (4.3.1)$$

explicitly has a non-positive slope. When calculating an inverse function the values  $C, N, m$  need to be chosen high enough in order to reach sufficiently high values for

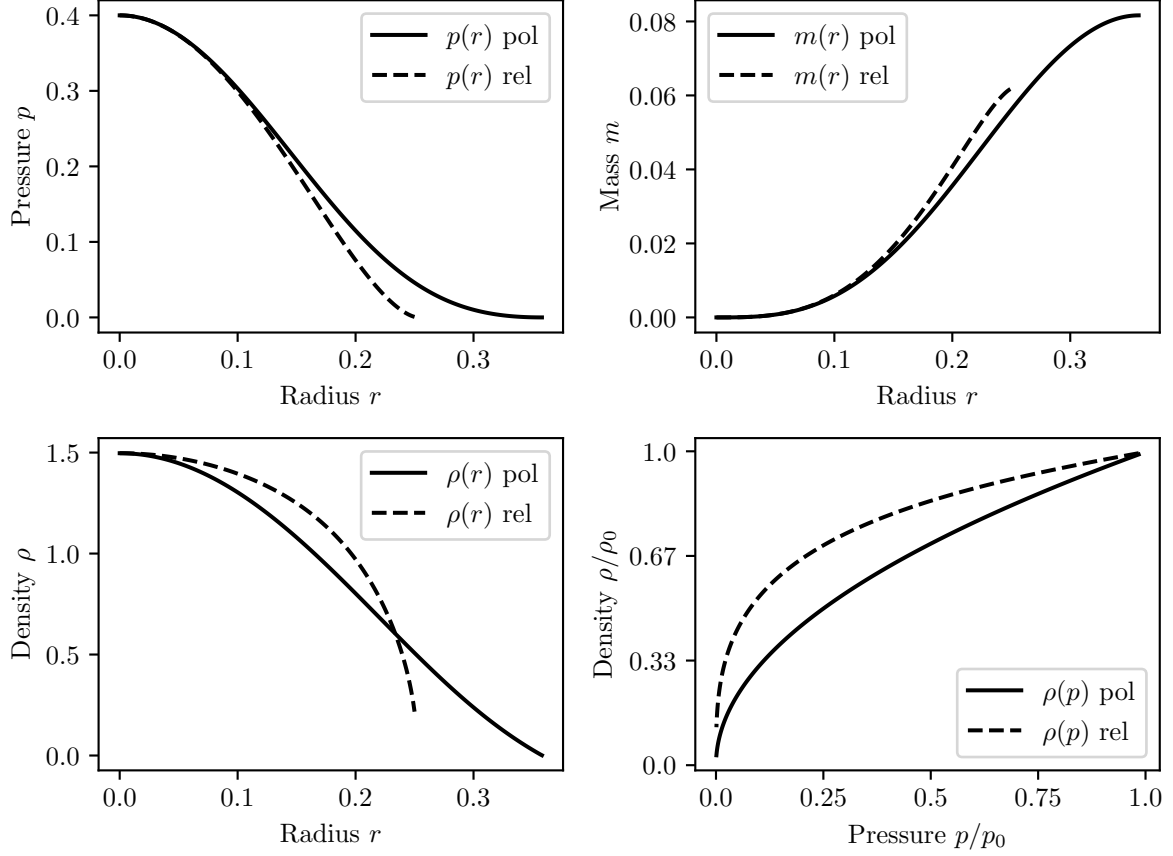


Figure 5: Comparison of TOV polytropic and relativistic EoS

$p$ . By knowing the initial value  $p_0$  and by the fact that the pressure of the TOV solution (3.1.22) has a non-positive slope and with equation (2.3.12), it is sufficient to choose  $Bp_0 = CNmc^2 > 2/e^2$ . On the other hand it is also important to choose high enough values for  $\alpha$  for numerical inversion since otherwise functions can not be displayed if they reach small enough values for  $p$ . The limit for  $\alpha$  is chosen such that values as low as  $0.001p_0$  can be evaluated. Afterwards the function over the given interval is inverted by a polynomial fit and then used in equation (2.3.10) to obtain the EoS.

Figure 5 is an example which shows a clear difference between the polytropic and relativistic EoS. The values used are  $B = 20$ ,  $p_0 = 0.4$ ,  $dr = 0.001$ ,  $n = 1$ . The correct prefactor  $A$  of the polytropic EoS is chosen after numerical calculation of the relativistic EoS. It is then simply given by

$$A = \frac{\rho_{\text{rel}}(p_0)}{p_0^{1/\gamma}}. \quad (4.3.2)$$

This normalises both EoSs to same initial pressures.

#### 4.4 Zero Values of the TOV and LE equation

Having discussed individual solutions for the TOV and LE equation, this section now turns to analysing zero values of both equations. We define a zero value as the first point in which the solution of our differential equation reaches value 0, or more precisely where

the pressure  $p$  reaches zero.<sup>3</sup> Even now the problem is not well posed since for different solving routines alternative transformations may be used. This has to be kept in mind when comparing results. Apart from the previously defined exact solutions of Table 1 we will only use the parameter  $r$  as defined by equation (3.1.24) to display final results. In order to solve the equations, a transformation in the form applied in section 3.2 is used since it was hoped to increase the precision with which the zero value is determined. This can be motivated by looking at plots for LE solutions as in Figure 2 and comparing them to TOV results of Figure 3. With the EoS  $\rho = A p^{1/\gamma}$  and the new definition  $\rho = \rho_0 \theta^n$ , the equations used are

$$\frac{\partial m}{\partial \xi} = 4\pi \kappa^3 \rho \xi^2 \quad (4.4.1)$$

$$\frac{\partial \theta}{\partial \xi} = - \frac{\left(1 + K \rho_0^{1/n} \theta\right) \left(4\pi \xi^3 \kappa^3 K \rho_0^{1+1/n} \theta^{n+1} + m\right)}{\left((n+1) K \rho_0^{1/n} \kappa \xi^2\right) \left(1 - \frac{2m}{\kappa \xi}\right)} \quad (4.4.2)$$

where  $\rho = \rho_0 \theta^n$  and initial values  $\partial m / \partial \xi = 0$  and  $\partial \theta / \partial \xi = 0$  can again be calculated with L'Hospitals rule together with  $\theta_0 = 1$  and  $m_0 = 0$ . When comparing equations (4.4.2) and (3.1.24) one can immediately recognise the distinct terms. When talking about zero values we stated earlier that transformations need be accounted when comparing results. However in this case since

$$p = \frac{\rho_0^n \theta^n}{A^{1+1/n}} \quad (4.4.3)$$

it is clear that zero values of  $\theta$  and  $p$  match under the same parameter  $\xi$ . When transforming back by using  $\kappa \xi = r$  as explained in section 3.2 one should note that  $\kappa$  depends on  $n$  via

$$4\pi \kappa^2 = (n+1) K \rho_0^{1/n-1} \quad (4.4.4)$$

and it will thus change values by more than constant scaling. Figure 6 shows zero values for solutions of the TOV and LE equations respectively with their dependence on  $n$ . Since not only the parameters  $A$  and  $n$  in the EoS may alter results but also the initial value  $p_0$  a family of curves was plotted. The upper half of the plot displays a currently not explainable bump which is prominent in the TOV  $p_0 = 0.1$  case. A similar but dampened behaviour can also be seen along the  $p_0 = 1$  curve. It is however still unclear if the TOV solutions similarly to the LE solutions blow up at the same combination of parameters. In order to further analyse this property higher integration ranges for larger values of the polytropic index  $n$  would be required. This observation motivates a theorem that for every combination of parameters  $A$  and  $p_0$  there exists a value  $n_0$  such that each solution of the TOV equation has no zero values for that particular combination of parameters  $A, p_0$  if  $n \geq n_0$ . The methods to prove this theorem will be discussed in the next section. To obtain Figure 6 the differential equations need to be solved multiple times for varying parameter combinations until the pressure reaches its zero value. Numerical optimisations to speed up calculation times, and a short discussion on the methods used can be found in the appendix in section D.

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<sup>3</sup>Note that  $p_0$  still denotes the initial value while  $r_0$  and  $\xi_0$  denote values where  $p = 0$ .

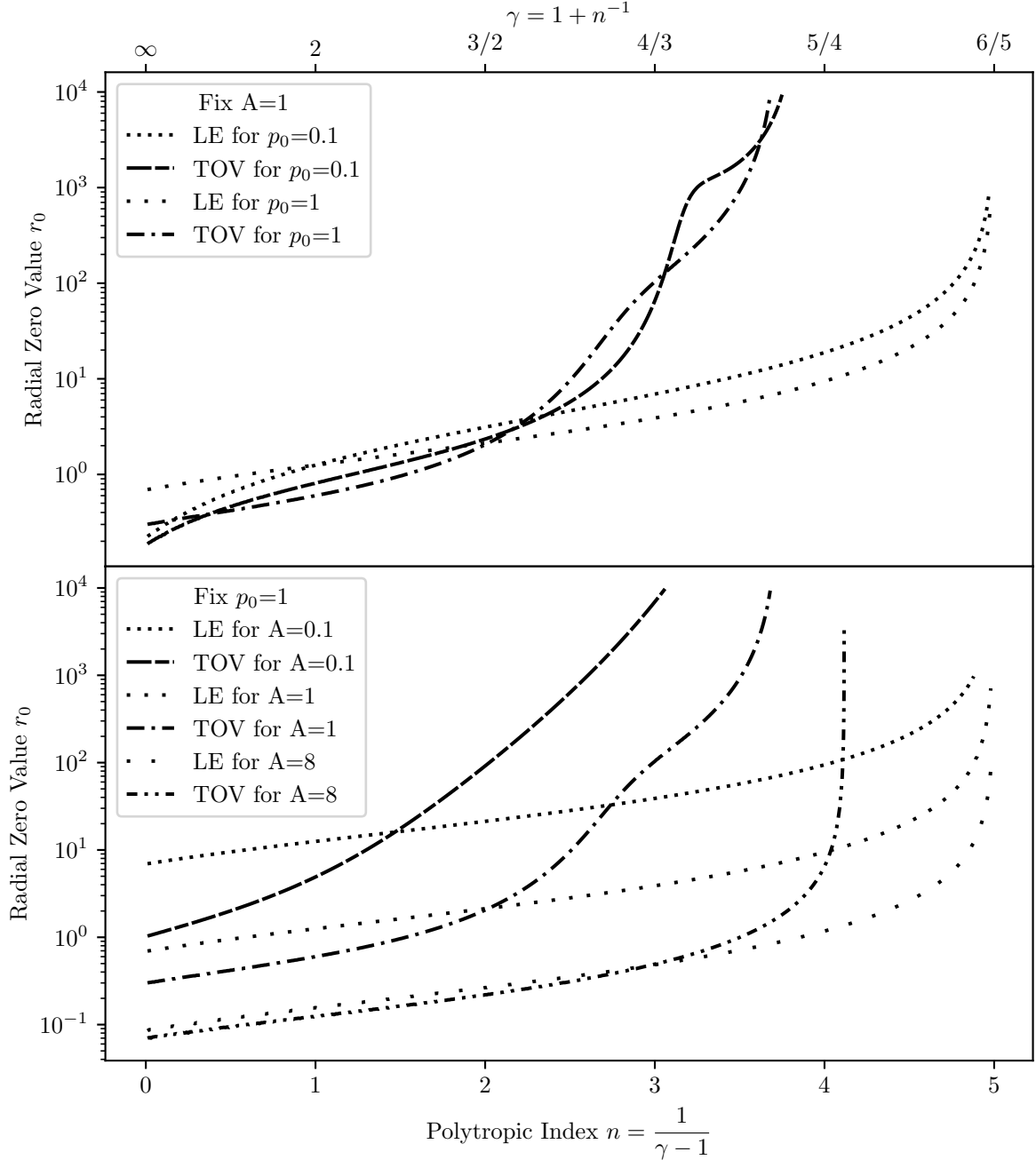


Figure 6: Zero Values of TOV and LE equation

In the first plot results for  $p_0 = 8$  have been omitted for better visual clarity. These behave similar to the  $p_0 = 1$  results.

## 5 To a General Zero Value Theorem

### 5.1 LE Equation

Motivated by the preceding sections, the aim now is to prove the following theorem.

**Theorem 5.1 - Lane-Emden Finite Boundary**

For every  $0 \leq n < 5$ , the generalised Lane-Emden equation

$$\frac{1}{\xi^{l-1}} \frac{d}{d\xi} \left( \xi^{l-1} \frac{d\theta}{d\xi} \right) + \theta^n = 0 \quad (5.1.1)$$

with  $\theta_0 = 1$  and  $l \geq 2$  has a zero value for a finite  $\xi_0$ . To prove this theorem, we show the statements following in this section.

**Lemma 5.2 - LE Local Existence**

For every  $n \geq 0$ , the Lane-Emden equation 5.1.1 with initial values defined as above has a unique solution in an  $\epsilon$  Ball around  $\xi = 0$ .

The procedure for proving existence is similar for both equations. Contrary to what is written in [QS07, p. 50] we do not expect to be able to solve this problem with the standard Banach fixed-point theorem. Instead, we want to use the Schauder fixed-point theorem [Sch30], which is a variation the Banach fixed-point theorem to show that over a suitably chosen space of functions we obtain a unique solution to the differential equation. In order to do this we examine some properties that solutions of the LE differential equation will have.

**Lemma 5.3**

A positive solution  $\theta$  of the LE equation with  $n \geq 0$  on a suitably chosen interval  $[0, \lambda]$  obeys the conditions

$$\theta(0) = \theta_0 > 0 \quad (5.1.2)$$

$$\left. \frac{d\theta}{d\xi} \right|_{\xi=0} = 0 \quad (5.1.3)$$

$$x \leq y \Rightarrow \theta(y) \leq \theta(x) \quad (5.1.4)$$

$$\theta_0/2 \leq \theta \quad (5.1.5)$$

$$|\theta(x) - \theta(y)| \leq L|x - y| \quad (5.1.6)$$

where  $L$  is independent of  $\theta$ . The statements above are true for every  $x, y \in [0, \lambda]$ . We will call  $C^* \subseteq C^0$  the space of all continuous functions  $f : [0, \lambda] \rightarrow \mathbb{R}$  that satisfy those conditions.

*Proof.* The first and second initial conditions of equations (5.1.2) and (5.1.3) are true by assumption. The third condition (5.1.4) can be motivated when once integrating the LE equation

$$\frac{d\theta}{d\xi} = - \int_0^\xi \frac{t^{l-1}}{\xi^{l-1}} \theta^n dt \quad (5.1.7)$$

where we not yet needed any of the above criteria except positivity of  $\theta$ . From this expression condition (5.1.6) is derived immediately. When integrating equation (5.1.7) once more and using the initial value of condition (5.1.2), we obtain

$$\theta = \theta_0 - \int_0^\xi \int_0^s \left(\frac{t}{s}\right)^{l-1} \theta^n dt ds. \quad (5.1.8)$$

We combine this statement with the previously shown condition of equation (5.1.4). Then it is clear that

$$\int_0^\xi \int_0^s \left(\frac{t}{s}\right)^{l-1} \theta^n dt ds \leq \theta_0^n \int_0^\xi \int_0^s \left(\frac{t}{s}\right)^{l-1} dt ds \leq \theta_0^n \frac{\xi^2}{2l} \quad (5.1.9)$$

and thus  $\xi \leq \lambda$  can be chosen suitably small to satisfy condition (5.1.5). For the last condition we again observe that from equation (5.1.8) follows promptly that

$$|\theta(x) - \theta(y)| = \left| \int_x^y \int_0^s \left(\frac{t}{s}\right)^{l-1} \theta^n dt ds \right| \leq \theta_0^n \frac{|x^2 - y^2|}{2l} \leq \theta_0^n \frac{\lambda}{l} |x - y| \quad (5.1.10)$$

□

The now introduced function space  $C^*$  contains all functions that are possible solutions for the LE equation. In order to use the Schauder fixed-point theorem, we have to show that  $C^*$  is compact and convex which we will show in the next two lemmata.

#### Lemma 5.4

The space  $C^*$  defined in lemma 5.3 with the norm  $\|\cdot\|_\infty$  is complete and convex.

*Proof.* We begin by showing completeness. It is clear that from condition (5.1.2) to (5.1.5) are conserved even in pointwise convergence. Generally speaking Lipschitz continuous functions need not be complete under  $\|\cdot\|_\infty$ . However, since we had chosen  $L$  to be independent of  $\theta$  for  $\|f_m - f\|_\infty \rightarrow 0$  we can write

$$|f(x) - f(y)| = |f(x) - f_m(x) + f_m(x) - f_m(y) + f_m(y) - f(y)| \leq 2\epsilon + L|x - y|$$

which proves the statement in the limit  $\epsilon \rightarrow 0$ . Convexity for conditions (5.1.2) to (5.1.5) is also clear. Let  $f, g \in C^*$  and  $h \in [0, 1]$ . Then

$$|hf(x) + (1-h)g(x) - hf(y) - (1-h)g(y)| \quad (5.1.11)$$

$$\leq h|f(x) - f(y)| + (1-h)|g(x) - g(y)| \quad (5.1.12)$$

$$\leq L|x - y| \quad (5.1.13)$$

which shows convexity. □

#### Lemma 5.5

The space  $C^*$  with  $\|\cdot\|_\infty$  is compact. To prove this statement we use the Arzela-Ascoli theorem. This means that since  $C^*$  is a function space on a compact set  $[0, \lambda]$  we only need to show that it is bounded and equicontinuous.

*Proof.* For every function  $f \in C^*$  we know by condition (5.1.4) and (5.1.5) that  $\theta/2 \leq f \leq \theta_0$  which shows boundedness. Since with condition (5.1.6) every function  $f \in C^*$  is especially Lipschitz, we can clearly see that equicontinuity holds.  $\square$

*Proof LE Local Existence.* With these results it is now possible to prove our initial statement of theorem 5.2. It is left to prove that the operator  $T : C^* \rightarrow C^*$  defined by

$$T(f)(x) = \theta_0 - \int_0^x \int_0^s \left(\frac{t}{s}\right)^{l-1} f^n(t) dt ds \quad (5.1.14)$$

is continuous on  $C^*$ . With a quick integral approximation we obtain

$$\|T(f) - T(g)\|_\infty = \left\| \int_0^x \int_0^s \left(\frac{t}{s}\right)^{l-1} [f^n(t) - g^n(t)] dt ds \right\|_\infty \quad (5.1.15)$$

$$\leq \frac{x^2}{2l} \|f^n - g^n\|_\infty. \quad (5.1.16)$$

Since  $z \mapsto z^n$  is continuous for  $z \geq 0$  and  $n \geq 0$  the statement follows directly.  $\square$

Equation (5.1.7) shows that non-negative solutions must have non-positive slope. The following theorem makes use of that fact and of a different representation of the LE equation to show that solutions can be extended until they reach zero.

### Lemma 5.6

Let  $\theta$  be a LE solution and  $[0, \xi_m]$  be the maximum interval of existence. Then  $\theta(\xi_m) = 0$  or  $\xi_m = \infty$ .

*Proof.* It is clear that  $\theta < 0$  is a case that can never appear as long as we are interested in positive solutions. Suppose  $\theta(\xi_m) > 0$  and  $\xi_m < \infty$ . Then there exists an interval  $(\xi_m - \epsilon, \xi_m + \epsilon)$  and  $(\theta(\xi_m) - \delta, \theta(\xi_m) + \delta)$  where the LE differential equation has a solution. With equations 4.1.7 we can directly see that the Picard-Lindelöf theorem [Lin94] is applicable in this situation.  $\square$

*Proof Lane-Emden Finite Boundary.* To prove Theorem 5.1 it is left to show that the LE equation has no global solutions for  $n < 5$ . In [QS07, p. 36], the author shows that for  $n \leq 5$ , LE solutions can not be global solutions. Furthermore it is explicitly shown that for  $n \geq 5$  global solutions exist and thus no zero-value exist. A variation of the proof for this particular usecase can be found in appendix E.  $\square$

## 5.2 TOV Equation

### Hypothesis 5.7 - TOV Zero Value Hypothesis

Given the TOV differential equation with  $\rho = A p^{\frac{n}{n+1}}$  and  $p_0, A > 0$

$$\frac{\partial m}{\partial r} = 4\pi \rho r^2 \quad (5.2.1)$$

$$\frac{\partial p}{\partial r} = -\frac{m\rho}{r^2} \left(1 + \frac{p}{\rho}\right) \left(\frac{4\pi r^3 p}{m} + 1\right) \left(1 - \frac{2m}{r}\right)^{-1} \quad (5.2.2)$$

and initial values  $m_0 = 0, p_0 \geq 0$  there exists a  $n_0 \geq 0$  such that all solutions with same parameters  $A, p_0$  and smaller exponent  $n < n_0$  have a  $p(r_0)$  for some  $r_0 > 0$ .



*Proof.*

□

**Lemma 5.8**

Let  $(m, p) : \mathbb{R} \rightarrow \mathbb{R}^2$  be a TOV solution with initial values  $m_0 = 0$  and  $p_0 > 0$ . Also let  $\rho : \mathbb{R}_{\geq 0} \rightarrow \mathbb{R}_{\geq 0}$  be the monotonously increasing EoS. Then for a small enough  $\epsilon$  Interval  $[0, \epsilon]$  the solution  $f$  satisfies the following conditions.

$$m(r) \leq \frac{4\pi\rho_0 r^3}{3} \quad (5.2.3)$$

$$\lim_{r \rightarrow 0} \frac{m}{r^3} = \frac{4\pi\rho_0}{3} \quad (5.2.4)$$

$$p(r) \geq \frac{p_0}{2} \quad (5.2.5)$$

*Proof.* To prove the condition (5.2.3) we inspect that with the initial value  $m_0 = 0$  we can write

$$m = \int_0^r 4\pi\rho r'^2 dr' \leq 4\pi\rho_0 \int_0^r r'^2 dr' = \frac{4\pi\rho_0 r^3}{3}$$

where in the second step we used that  $\rho$  is monotonously increasing and positive solutions of the TOV equation have non-positive slope in the pressure  $p$  component.

The second condition is obtained when applying L'Hospital's Rule as was done in the beginning of section 4.1. For the third condition we first inspect that condition 5.2.3 lets us choose  $(1 - 2m/r)^{-1} \leq 2$ .

$$p_0 - p = \int_0^r \frac{1}{r'^2} (p + \rho) (4\pi p r'^3 + m) \left(1 - \frac{2m}{r'}\right)^{-1} dr' \quad (5.2.6)$$

$$\leq 2(p_0 + \rho_0) \int_0^r \frac{4\pi p r'^3 + m}{r'^2} dr' \quad (5.2.7)$$

$$\leq 2(p_0 + \rho_0) \int_0^r \left(4\pi p_0 r' + \frac{4\pi\rho_0 r'}{3}\right) dr' \quad (5.2.8)$$

$$= 4\pi(p_0 + \rho_0) \left(p_0 + \frac{\rho_0}{3}\right) r^2 \quad (5.2.9)$$

In the first step we used that  $p$  is a TOV solution while in the second step the aforementioned inequality and just as in the last step the non-positive slope of  $p$  with the monotonicity of  $\rho$  came into play. Afterwards we used the first condition 5.2.3 and integrated the result. It is now clear that a small enough  $\epsilon > 0$  can be chosen such that all stated conditions are satisfied. □

**Lemma 5.9 - TOV Local existence**

The TOV equation 5.2.2 with initial values  $m_0 = 0$  and  $p_0 \geq 0$  and monotonously increasing continuous EoS  $\rho : \mathbb{R}_{\geq 0} \rightarrow \mathbb{R}_{\geq 0}$  has a local solution for  $r \in [0, \epsilon]$  with some  $\epsilon > 0$ .

*Proof.* The proof follows similar principles as previously done in lemma 5.2. In order to solve the problem at hand we need to define a space  $K$  consisting of all functions

$f : [0, \epsilon] \rightarrow \mathbb{R}_{\geq 0}^2$  that all satisfy the same boundary conditions in form of initial values  $m_0 = 0$  and  $p_0 > 0$ . Additionally, we have shown in lemma 5.8 that we can apply the listed conditions for a small enough interval  $[0, \epsilon]$  without reducing the number of possible solutions. Since solutions of the TOV equation have monotonously decreasing pressure  $p$  and monotonously increasing mass  $m$  those are taken as an additional requirement. We now call  $K$  the space of all functions  $f : [0, \epsilon] \rightarrow \mathbb{R}_{\geq 0}^2$  that satisfy these conditions. This space is complete with the norm  $\|\cdot\|_\infty$ . Let  $T : \tilde{K} \rightarrow K$  be defined by the TOV equation (3.1.22)

$$T((m, p))(r) = \left( \int_0^r 4\pi \rho r'^2 dr', - \int_0^r \frac{p + \rho}{r'^2} (4\pi \rho r'^3 + m) \left(1 - \frac{2m}{r'}\right)^{-1} dr' \right).$$

We want to prove that the integrands are Lipschitz in  $m, p$ . For the first component of  $T$  this statement is clearly true since  $\rho$  is monotonously increasing and  $p \in [p_0/2, p_0]$  is bound by assumption (5.2.5). For the second integrand together with restriction (5.2.3) we observe that

$$(p + \rho) \frac{4\pi \rho r'^3 + m}{r^2} \left(1 - \frac{2m}{r}\right)^{-1} \leq 2(p_0 + \rho_0) \left(4\pi \rho r' + \frac{m}{r}\right).$$

The first resulting term is again Lipschitz in  $p$  by the same argument. For the second term we can use restriction (5.2.4). For the second term we look at  $(m, p), (u, q) \in K$  with  $0 < \delta < 1$  Then for  $r \geq \delta$  it is clear that

$$\left\| \frac{m - u}{r^2} \right\|_\infty \leq \frac{1}{\delta^2} \|m - u\|_\infty$$

At  $r = 0$  the expression clearly vanishes due to constraint (5.2.3). Here any Lipschitz constant can be chosen. For  $0 < r < \delta$  we readily see that

$$\left\| \frac{m - u}{r^2} \right\|_\infty \leq \frac{1}{\delta^2} \|m - u\|_\infty$$

which finishes the proof. □

## 6 Ideas

### Theorem 6.1 - Lane-Emden Finite Boundary 2

For  $n < 5$ , the Lane-Emden equation 5.1.1 has no Solution defined on the total space  $\mathbb{R}_{\geq 0}$ .

*Proof.* From [QS07] we know that solutions for equation 5.1.1 on the total space exist for  $n \geq 5$ . □

### Theorem 6.2 - Lane-Emden Finite Boundary 3

Let  $n < 5$  and  $\theta : [0, x) \rightarrow \mathbb{R}$  be a solution of 5.1.1. If  $\theta(x) > 0$ , then  $\theta$  can be extended with Theorem 5.2 to  $[0, x + \epsilon)$  where  $\epsilon$  follows a growth condition  $\epsilon(\xi) = \dots$  such that if  $\theta$  would have no zero point, there would be a solution on total  $\mathbb{R}_{\geq 0}$  and thus a contradiction with 6.1.

### Theorem 6.3 - Lane-Emden Finite Boundary 4

The Lane Emden equation has a zero value if the exponent function  $n(\xi)$  fullfills the growth condition XYZ.

### Theorem 6.4 - TOV Finite Boundary

For each  $A > 0$ , the TOV equation 4.1.1 has a exponent  $n > 0$  for which a solution  $p$  does not have any zero points.

### Theorem 6.5 - TOV Finite Boundary 1

If  $p_1$  and  $p_2$  are two solutions of the TOV equation (for equal  $A$ ) with  $p_i$  corresponding to  $n_i$  exponents, then if  $n_1 < n_2$ , the zero point (if it exists) of  $p_1$  is smaller than that of  $p_2$  (if it exists or is  $\infty$ ).

### Theorem 6.6 - TOV Finite Boundary 2

a

## 7 Outlook

# Appendix

## A Known Exact Solutions of the LE Equation

This section relies on information in [Wei20] and [Cha58]. The LE equation is

$$\frac{1}{\xi^2} \frac{d}{d\xi} \left( \xi^2 \frac{d\theta}{d\xi} \right) + \theta^n = 0 \quad (\text{A.1})$$

which for  $n = 0$  transforms readily into

$$\int_0^\xi \frac{d}{d\xi} \left( \xi'^2 \frac{d\theta}{d\xi} \right) d\xi' = - \int_0^\xi \xi'^2 d\xi'. \quad (\text{A.2})$$

Both sides can be evaluated directly and then simplified further.

$$\xi^2 \frac{d\theta}{d\xi} = -\frac{\xi^3}{3} \quad (\text{A.3})$$

$$\theta(\xi) = \theta(0) - \frac{\xi^2}{6} \quad (\text{A.4})$$

With the initial condition  $\theta(0) = 1$ , we obtain the desired result. For  $n = 1$ , equation A.1 transforms into

$$\frac{d}{d\xi} \left( \xi^2 \frac{d\theta}{d\xi} \right) + \xi^2 \theta = 0 \quad (\text{A.5})$$

we start by substituting  $U = \theta x$  and thus obtain (also multiplying by  $x^2$ )

$$0 = \frac{dU}{d\xi} + \xi \frac{d^2 U}{d\xi^2} - \frac{dU}{d\xi} + \xi U \quad (\text{A.6})$$

$$-U = \frac{d^2 U}{d\xi^2} \quad (\text{A.7})$$

and the last equation can be solved with a linear combination of  $\cos(\xi)$  and  $\sin(\xi)$ . Transforming back to  $\theta$ , we then have

$$\theta(\xi) = A \frac{\sin(\xi)}{\xi} - B \frac{\cos(k\xi)}{\xi}. \quad (\text{A.8})$$

The need for a well defined limit at  $\xi \rightarrow 0$  implies that  $B = 0$  and thus since  $\sin(z)/z \rightarrow 1$  for  $z \rightarrow 0$ , we have  $A = \theta(0) = 1$  and

$$\theta(\xi) = \frac{\sin(\xi)}{\xi}. \quad (\text{A.9})$$

For  $n = 5$ , we start by making the two substitutions  $x = 1/\xi$  and  $\theta = ax^\omega$  as done in [Cha58, pp. 94 sqq.] and then transforming the LE equation into

$$x^4 \frac{d\theta}{d\xi} + \theta^n = 0 \quad (\text{A.10})$$

$$a\omega(\omega - 1)x^{\omega+2} + a^n x^{n\omega} = 0. \quad (\text{A.11})$$

From this we see that  $\omega + 2 = n\omega$  and  $a\omega(\omega - 1) = -a^n$  needs to be satisfied, since the equation needs to hold for all  $x \in \mathbb{R}_{\geq 0}$ . Rewriting these conditions, we obtain the singular solution for the LE equation

$$\theta(x) = \left( \frac{2(n-3)}{(n-1)^2} \right)^{1/(n-1)} x^{2/(n-1)} \quad (\text{A.12})$$

since with  $x = 1/\xi$ , we have  $\theta(\xi) \rightarrow \infty$  for  $\xi \rightarrow 0$ . Notice that the solution is only valid if  $n \geq 3$ . We can use this solution to the LE equation and perturb it to make the more general ansatz

$$\theta(x) = ax^\omega z(x). \quad (\text{A.13})$$

If  $n < 3$ , the factor  $a$  has to be replaced with a more general one. The singular solution is obtained when taking  $z = 1$ . Using another transformation  $1/x = \xi = \exp(-t)$ , we obtain

$$\frac{1}{\xi^2} \frac{d}{d\xi} \left( \xi^2 \frac{d\theta}{d\xi} \right) + \theta^n = 0 \quad (\text{A.14})$$

$$x^4 \frac{d^2\theta}{dx^2} + \theta^n = 0 \quad (\text{A.15})$$

$$ax^{\omega+2} \left[ x^2 \frac{d^2z}{dx^2} + 2\omega x \frac{dz}{dx} + \omega(\omega-1)z \right] + a^n x^{n\omega} z^n = 0 \quad (\text{A.16})$$

$$\frac{d^2z}{dt^2} + (2\omega-1) \frac{dz}{dt} + \omega(\omega-1)z + a^{n-1} z^n = 0 \quad (\text{A.17})$$

$$\frac{d^2z}{dt^2} + \frac{5-n}{n-1} \frac{dz}{dt} + 2 \frac{3-n}{(n-1)^2} z + 2 \frac{(n-3)}{(n-1)^2} z^n = 0. \quad (\text{A.18})$$

For  $n = 5$ , we obtain

$$\frac{d^2z}{dt^2} = \frac{1}{4} z(1 - z^4). \quad (\text{A.19})$$

we multiply both sides with  $dz/dt$  and integrate

$$\frac{1}{2} \left( \frac{d^2z}{dt^2} \right) = \frac{1}{8} z^2 - \frac{1}{24} z^6 + D \quad (\text{A.20})$$

where  $D$  is the integration constant. For  $\xi \rightarrow 0$ , we expect  $\theta \rightarrow \theta_0$  and thus  $z = \theta_0 e^{-\omega t} (1/a + \mathcal{O}(e^{-t}))$  as  $t$  approaches  $\infty$ . We immediately see that  $dz/dt$  exhibits a similar behaviour and thus the integration constant  $D$  has to vanish.

The right-hand side of equation A.20 cannot get negative since otherwise  $z$  would take complex values which enables us to take the square root with a minus sign<sup>4</sup> and integrate again

$$\int \left( 1 - \frac{1}{3} z^4 \right)^{-1/2} \frac{dz}{z} = -\frac{1}{2} \int dt. \quad (\text{A.21})$$

We change the integration by again substituting the variables  $1/3z^4 = \sin^2(\alpha)$  and calculate  $dz/z = 2 \cos(\alpha)/\sin(\alpha) d\alpha$  as well as  $-dt = d\alpha/\sin(\alpha)$ . Now we rewrite the integral

$$\int \frac{1}{\sqrt{1 - \sin^2(\alpha)}} \frac{2 \cos(\alpha)}{\sin(\alpha)} d\alpha = - \int dt \quad (\text{A.22})$$

$$\int \frac{2 d\alpha}{\sin(\alpha)} = - \int dt. \quad (\text{A.23})$$

---

<sup>4</sup>This only determines the direction in which  $t$  is defined, so it is arbitrary.

Evaluating those integrals leads us to

$$\log(\tan(\alpha/2)) + \log(1/C) = -t \quad (\text{A.24})$$

where the integration constant has been chosen in advance to simplify the next expressions. From here, we further manipulate equation A.24 and combine it with our previous substitution  $1/3z^4 = \sin^2(\alpha)$  to obtain

$$\frac{1}{3}z^4 = \sin^2(\alpha) = \frac{4 \tan^2(\alpha/2)}{(1 + \tan^2(\alpha/2))^2} \quad (\text{A.25})$$

and with the solution for our integral before and plugging in the substitution from the beginning  $\xi = e^{-t}$ , we have

$$z = \pm \left( \frac{12C^2\xi^2}{1 + C^2\xi^2} \right)^{1/4}. \quad (\text{A.26})$$

For  $\theta$ , we need to have  $\theta \rightarrow \theta_0 = 1$  as  $\xi \rightarrow 0$ . This means that  $C = 1$  and with  $\theta = ax^\omega z$ , we obtain

$$\theta = \frac{1}{(1 + \frac{1}{3}\xi^2)^{1/2}}. \quad (\text{A.27})$$

We see that this equation has no zero value and tends to 0 as  $\xi \rightarrow \infty$ .

## B New Exact LE Series Solution at Index 2

One can derive another exact solution for the LE equation at the exponent  $n = 2$  when considering a series expansion of  $\theta$  around the point  $\xi = 0$  with initial conditions

$$\theta = \sum_{m=0}^{\infty} a_m \xi^m \quad a_0 = \theta|_{\xi=0} = \theta_0 \quad a_1 = \left. \frac{d\theta}{d\xi} \right|_{\xi=0} = 0 \quad (\text{B.1})$$

Since the series is absolutely convergent, we can plug it into the LE equation 3.2.5 and use the Cauchy Product formula.

$$\sum_{m=2}^{\infty} m(m-1)a_m \xi^{m-2} + \sum_{m=1}^{\infty} (2m)a_m \xi^{m-2} + \sum_{m=0}^{\infty} \sum_{k=0}^m a_{m-k} a_k \xi^m = 0 \quad (\text{B.2})$$

### Theorem B.1

The odd coefficients  $a_{2m+1}$  of this series expansion vanish.

*Proof.* We rewrite the summations of equation B.2 to start at the same index  $m = 0$  and combine them

$$\sum_{m=0}^{\infty} \left( (m+2)(m+1)a_{m+2}\xi^m + (2m+2)a_{m+1}\xi^{m-1} + \sum_{k=0}^m a_{m-k}a_k\xi^m \right) = 0 \quad (\text{B.3})$$

With equation B.1, we can start the summation in the middle one index higher and separate the term  $\xi^m$ . This equation has to be true inside the radius of convergence of the series and thus needs to vanish for ambiguous  $\xi$ .

$$(m+2)(m+1)a_{m+2} + 2(m+2)a_{m+1} + \sum_{k=0}^m a_{m-k}a_k = 0 \quad (\text{B.4})$$

and upon further manipulation results in the recursive description for the coefficients of the series

$$a_{m+2} = -\frac{1}{(m+2)(m+3)} \sum_{k=0}^m a_{m-k} a_k. \quad (\text{B.5})$$

We show the statement by induction. For  $a_1$  it is already true. Let the statement be true for all odd values  $2k+1 \leq 2m+1$ . Writing down  $a_{2m+3}$  gives us

$$a_{2m+3} = -\frac{1}{(2m+3)(2m+4)} (a_0 a_{2m+1} + a_1 a_{2m} + \cdots + a_{2m} a_1 + a_{2m+1} a_0). \quad (\text{B.6})$$

It is clear that in the summation odd and even coefficients get paired and will thus vanish completely.  $\square$

This proof shows that we can restrict ourselves to the subsequence  $b_m = a_{2m}$  and the subseries with initial values given by

$$\theta = \sum_{m=0}^{\infty} b_m \xi^{2m} \quad b_{m+1} = -\frac{1}{(2m+2)(2m+3)} \sum_{k=0}^m b_{m-k} b_k \quad b_0 = \theta_0 \quad (\text{B.7})$$

### Theorem B.2

The series  $\theta = \sum_{m=0}^{\infty} b_m \xi^{2m}$  converges for  $\xi \leq 1$ .

*Proof.* We start by showing that  $|b_{m+1}| \leq 1/(4m+6)$ . Using the triangle inequality on B.7, we obtain

$$|b_{m+1}| \leq \frac{1}{(2m+2)(2m+3)} \sum_{k=0}^m |b_{m-k} b_k|. \quad (\text{B.8})$$

For  $m=0$ , we have  $b_1 = -1/6$  and thus  $|b_1| \leq 1/3$  and so the statement is true for  $m=0$ . If the statement holds for all  $m < m+1$ , we have

$$|b_{m+1}| \leq \frac{m}{(2m+2)(2m+3)} = \frac{1}{2} \frac{m}{(m+1)} \frac{1}{(2m+3)} \leq \frac{1}{2} \frac{1}{(2m+3)} \quad (\text{B.9})$$

which proves the first statement. To see that  $b_m$  is an alternating sequence, we again inspect B.7. The first element of the series is  $b_0 = 1$  and the second is  $b_1 = -1/6$  so the initial statement is again correct. Let  $b_m$  be alternating for all  $m < 2m+1$ . Then equation B.7 for odd values reads

$$b_{2m+1} = -\frac{1}{(4m+4)(4m+5)} (b_{2m} b_0 + b_{2m-1} b_1 + \cdots + b_0 b_{2m}) \quad (\text{B.10})$$

This shows that if all odd values  $b_{2k+1}$  are negative and all even ones are positive, then  $b_{2m+1}$  is negative too. The same holds true when the index  $2m+2$  is even. By the Leibniz criterion the series converges if  $\xi \leq 1$ .  $\square$



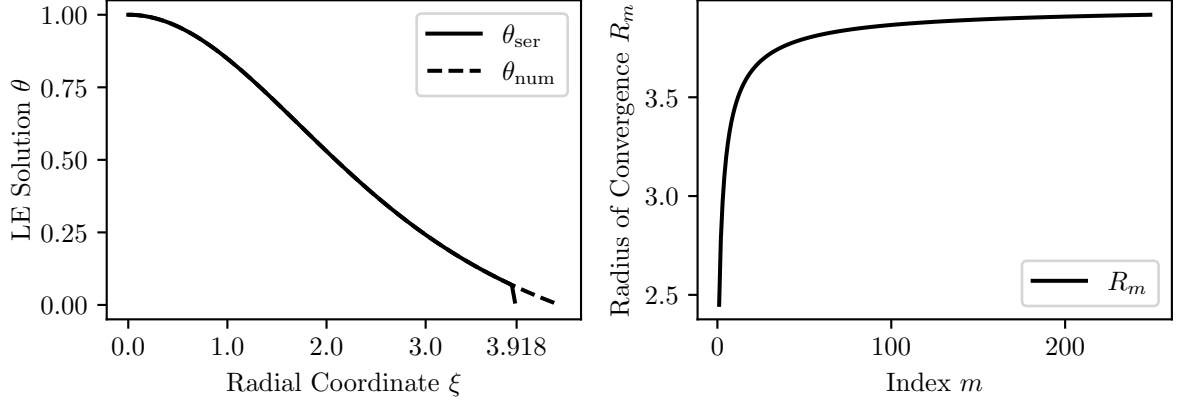


Figure 7: LE Solution for  $n = 2$

We see the power the series calculated with the coefficients explained above. The far right tick is the last calculated value for  $R_m$  where plotting is stopped) and the sequence  $R_m = (|a_m|)^{-1/m}$  which approaches the radius of convergence.

These results can be used to calculate the power series of the LE equation numerically. Figure 7 shows those results. Good agreement is found between the two solutions as the difference  $\Delta = \theta_{\text{calc}} - \theta_{\text{ser}}$  is not larger than 0.00015. The solutions differ for larger values around  $\approx 3.918$ . The top right plot indicates that the radius of convergence might actually be larger than where solutions deviate. One important observation is that in this procedure, the coefficients were calculated up to  $b_{250} \approx -4.7723 \times 10^{-296}$  which is in proximity to the floating point limit of python given by  $\approx 2.2251 \times 10^{-308}$  and indicates that numerical errors may be the reason the series starts to diverge. Also in order to achieve sufficient numerical precision, we would need to calculate higher order terms of the series, which is supported by  $(R_{250})^{500} \approx 3.2159 \times 10^{296}$  and shows that these parts of the series still yield large contributions to the value of  $\theta$  at this point.

## C Exact TOV Solution in Absence of Mass

### Theorem C.3 - TOV Exact Solution

The TOV equation 4.1.1 with a polytropic EOS  $\rho = Ap^{1/\gamma}$  has a well defined limiting case where  $A \rightarrow 0$  with  $m = 0$  and

$$p = \frac{p_0}{2\pi r p_0 + 1} \quad (\text{C.1})$$

*Proof.* First we transform the TOV equation 3.1.24 using  $p = y^a$  with  $a > 0$  and  $m = Arv$  and together with the polytropic EOS  $\rho = Ap^{1/\gamma}$  obtain

$$\frac{\partial v}{\partial r} = 4\pi r y^{a/\gamma} - Av \quad (\text{C.2})$$

$$ay^{a-1} \frac{\partial y}{\partial r} = -\frac{vA^2 y^{a/\gamma}}{r} \left(1 + \frac{y^a}{Ay^{a/\gamma}}\right) \left(\frac{4\pi r^2 y^a}{Av} + 1\right) \frac{1}{1 - 2vA} \quad (\text{C.3})$$

Rearranging the second equation one obtains

$$\frac{\partial y}{\partial r} = -\frac{y^{a/\gamma - a + 1}}{ar} (A + y^{a - a/\gamma}) (4\pi r^2 y^a + Av) \frac{1}{1 - 2vA}. \quad (\text{C.4})$$

Using  $\gamma = 1 + 1/n$  with  $n > 0$ , we see that this equation is continuous in every variable  $(r, y, v) \in \mathbb{R}_{>0} \times \mathbb{R}_{\geq 0} \times [0, 1/2A)$ . We restrict ourselves to a compact domain  $r \in \overline{B_\delta(\tau)}$  and  $v \in [0, 1/2A - \epsilon]$  where  $0 < \epsilon < 1/2A$ . We choose  $0 < \delta$  and  $0 \leq \tau$  in such a way that  $0 < \epsilon < 1/2A$  is satisfied for some  $\epsilon$  which is possible since  $m \rightarrow 0$  for  $r \rightarrow 0$ . To obtain Lipschitz continuity in  $(y, v)$ , all of the following conditions need to be fulfilled.

$$\frac{a}{\gamma} - a + 1 \geq 1 \quad a - \frac{a}{\gamma} \geq 1 \quad a \geq 1 \quad (\text{C.5})$$

The second equation implies the first and the third. Thus we only need to choose  $a \geq (1 - 1/\gamma)^{-1}$ . With  $\gamma = 1 + 1/n$  we can rewrite this equation to

$$a \geq n + 1 \quad (\text{C.6})$$

which can be easily satisfied. This now shows with extension of the Picard-Lindelöf Theorem that there exists a unique solution for given initial values  $\tau, y(\tau), v(\tau) \in B_\delta(\tau) \times \mathbb{R}_{\geq 0} \times [0, 1/2A - \epsilon]$  that especially continuously depends on  $A$ . Without loss of generality we can choose  $\delta$  small enough such that solutions are positive in  $p$ . By transforming back equation C.3 to

$$\frac{\partial p}{\partial r} = -\frac{1}{r^2} (Ap^{1/\gamma} + p) (4\pi r^3 p + vA) (1 - 2vA)^{-1} \quad (\text{C.7})$$

and letting  $A \rightarrow 0$ , we obtain

$$\frac{\partial p}{\partial r} = -4\pi r p^2 \quad (\text{C.8})$$

which is then solved by

$$p = \frac{\tilde{p}}{2\pi\tilde{p}(r^2 - \tau^2) + 1} \quad (\text{C.9})$$

where  $\tilde{p}$  is the initial value at  $r = \tau$ . It is clear that this solution can be extended to  $r \in [0, \infty)$  by choosing arbitrary small  $\tau$ . In the limit, the result equals the hypothesis.  $\square$

## D Numerical Optimisations

When solving the TOV or LE equation numerically calculations for the next step clearly depend on results from the before calculated one. This simple fact which is integral to the concept of an ODE means that numerics is single thread <sup>5</sup> bound which already puts one at a systematic disadvantage. This means any effort to parallelise the given task must focus on distributing the individual solving routines to respective threads. With this first consideration at hand a server with high number of processing units may come into mind to then speed up the calculation. However in the situation of section 4.4 this would only allow us to calculate more results in parallel but not speed up the individual solving routines. And this is exactly the bottleneck in our case. When looking at figure 6 it is clear that higher values for  $r_0$  are of interest where the curves reach the current maximum limit. From this figure it is also clear that the effort will increase exponentially since we are bound to a single thread for individual solving routines. Hence parallel computation of this problem would not help in this case especially considering that server processing units typically have lower single threaded performance than consumer chips.

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<sup>5</sup>Thread refers to a sequence of instructions given to a processing unit.

Explicit parallelisation is achieved by distributing different values of  $n$  to different threads. More efficient distribution of those calculations can make a difference but effects discussed above and below have more impact. Furthermore in order to fully understand how to improve for this distribution it is partly necessary to calculate the results which is then even less of an issue afterwards.

Another optimisation evolved around not computing unnecessary cases. Consider again figure 6. It is clear (however not proven until now) that for each combination  $A, p_0$  at some value  $n_0(A, p_0)$  values for  $r_0$  will reach the numerical upper limit of the routine. For other solving routines with  $n \geq n_0$  the same will happen which makes it unnecessary to even begin to compute them since they will not be shown in the graph anyways. This behaviour has been verified for a number of numerical combinations of parameters before being utilised in the final version. For the solving routine this means explicitly that once there have been 2 previous solving routines that failed in obtaining a value lower than  $r_0$ , the whole routine will not consider solving for any combinations of these particular parameters  $A, p_0$  anymore.

Another obvious optimisation was to use a database to store results. In this case MongoDB [DH21] was chosen for its good integration with the python language. This pool of results allows one to compare the current solving routine with already calculated results. Extra calculations can be permitted after a database query in the following cases:

- If the current solving routine does have higher precision (in form of lower constant stepsize).
- If the previous solving routine reached the limit for  $r_0$  and the current one would go further.
- If there exist multiple results for a combination of parameters  $A, p_0, n$  <sup>6</sup>.

With these optimisations less results in total need to be calculated. One still present disadvantage is however that allocation of parameters for solving routines is done beforehand and can not be changed in mid process. This means if a thread started by the program finished and another is still running with multiple solving routines left, these cannot be given to another thread. This is a possibility for future tweaks.

So far the previous discussion evolved around a fixed stepsize. Since in this case the 4th order Runge Kutta solving method was implemented by hand, it is possible to use variable stepsizes for solving. From figure 4 we can see that the largest errors occur in the first step of integration. This means in the beginning a small stepsize is desirable to retain accuracy while later it can be increased. These settings were experimented with but since the author had already achieved numerous results with a constant stepsize were not committed to the final version. It is however believed to yield a significant advantage if one would want to calculate higher values of  $r_0$ .

Furthermore, one could attempt to rewrite the code in a lower level programming language such as C to obtain a performance benefit. Results from [Jen21] show a 676% increase in performance for calculating Fibonacci numbers between the two languages when compiling code with PyPy. It is however unclear to which degree this theoretical improvement would carry over to the code present.

The code is also designed in a way to check regularly for already present results in the

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<sup>6</sup>Then every result present is deleted and one new result will be calculated.

database. These calls could be forwarded to a different handler in order to possibly parallelise the problem and reduce disk<sup>7</sup> latency issues.

## **E Nonexistence of Global LE Solutions**

We present a variation of the proof given in [QS07, p. 36] to conclusively show theorem 5.1.

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<sup>7</sup>We refer to the term disk although the data was stored on a solid state drive.

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