New Solutions of Tolman-Oppenheimer-Volkov-Equation and of Kerr Spacetime with Matter and the Corresponding Star Models

Jan Helm ©

Technical University of Berlin, Berlin, Germany
Email: jan.helm@alumni.tu-berlin.de

Abstract

The Tolman-Oppenheimer-Volkov (TOV) equation is solved with a new ansatz: the external boundary condition with mass $M_0$ and radius $R_1$ is dual to the internal boundary condition with density $\rho_{bc}$ and inner radius $r_i$, and the two boundary conditions yield the same result. The inner boundary condition is imposed with a density $\rho_{bc}$ and an inner radius $r_i$ which is zero for the compact neutron stars, but non-zero for the shell-stars: stellar shell-star and galactic (supermassive) shell-star. Parametric solutions are calculated for neutron stars, stellar shell-stars, and galactic shell-stars. From the results, an M-R-relation and mass limits for these star models can be extracted. A new method is found for solving the Einstein equations for Kerr space-time with matter (extended Kerr space-time), i.e. rotating matter distribution in its own gravitational field. Then numerical solutions are calculated for several astrophysical models: white dwarf, neutron star, stellar shell-star, and galactic shell-star. The results are that shell-star star models closely resemble the behaviour of abstract black holes, including the Bekenstein-Hawking entropy, but have finite redshifts and escape velocity $v < c$ and no singularity.

Keywords

General Relativity, Tolman-Oppenheimer-Volkov Equation, Neutron Stars, Shell Stars

1. Introduction

In General Relativity, one of the most important applications is to calculate the mass distribution and the space-time metric for a given equation-of-state of a stellar model.
Without rotation, one has spherical symmetry and then the Tolman-Oppenheimer-Volkov (TOV) equation in radius $r$, which is derived directly from the Einstein equations (see [1]), is being used. The TOV equation consists originally of 2 coupled non-linear ordinary differential equations (odeq) of degree 1 in $r$ for mass $M(r)$ and density $\rho(r)$, where $4\pi r^2 \rho(r) = M'(r)$, and can be transformed into one ordinary differential equation of degree 2 for $M(r)$ by eliminating $\rho(r)$.

The boundary condition is imposed normally at $r = 0$ with $M(0) = 0$ and $\rho(0) = \rho_0$, where $\rho_0$ is the maximal density. Then the TOV-equation for $M(r)$ is solved with this boundary condition at $r = 0$ for $M(r)$ and $M'(r)$, which gives the total mass $M_0(\rho_0)$ and the total radius $R(\rho_0)$, and a mass-radius relation $M_0(R)$.

The predominant view of the neutron stars and stellar black-holes is, that neutron stars obey an equation-of-state (eos) of an interacting-fluid model [2], which solutions of the TOV equation up to about $M = 3M_{\odot}$. For larger masses, it is assumed that only a black-hole solution remains. This is based on the so-called Oppenheimer limit for the radius of a compact mass

$$R_{\text{lim}} = \frac{9}{8} r_i = \frac{9MG}{4c^2};$$

for a smaller compact object, the density at the center becomes infinite.

The new ansatz presented here is the extended (inner) boundary condition at $r = r_i$ with the non-zero inner radius $r_i$, $M(r_i) = 0$ and $\rho(r_i) = \rho_0$, i.e. the star becomes a shell-star with an (almost) void interior. With the parameters $r_i$ and $\rho_0$, this ansatz generates a 2-parametric solution manifold, where, because of energy minimization, the stable physical solution is the one with minimal $r_i$ for a given $\rho_0$, which determines the total mass $M_0(r_i, \rho_0)$ and the total radius $R(r_i, \rho_0)$. This ansatz circumvents the Oppenheimer limit, because the mass is non-compact, and the density at the center is zero. It yields valid solutions of the TOV-equation with a continuous mass-radius relation $M_0(R, r_i)$.

The Oppenheimer limit $R_{\text{lim}} = \frac{9}{8} r_i = \frac{9MG}{4c^2}$ is deduced from the maximum pressure at $r = 0$:

$$P(0) = \rho_0 c^2 \frac{1 - \sqrt{1 - \frac{r_i}{R}}}{3 \sqrt{1 - \frac{r_i}{R} - 1}}$$

The dual (outer) boundary condition is the one at $r = R$ with $M = M_c$ and $\rho = \rho_c$, where $\rho_c$ depends on the equation-of-state (eos): for neutron stars with interacting nucleon fluid with the equilibrium nucleon density $\rho_c$, and for nucleon Fermi-gas (stellar shell-stars) $\rho_c = 0$.

The 2 parameters $R$ and $M_c$ in the dual outer boundary condition correspond uniquely to the 2 parameters $r_i$ and $\rho_0$ in the inner boundary condition.

With rotation, one has an axisymmetric model in the variables $r$ and $\theta$ (azimuthal angle), and has to solve the Einstein equations in these 2 coordinates. In
vacuum, the corresponding solution is the Kerr space-time in $r$ and $\theta$.

With mass, a good starting point is using the extended Kerr space-time in Boyer-Lindquist coordinates with correction-factor functions $A_0, \ldots, A_4$ and $B_0, \ldots, B_4$ and the mass $M(r, \theta)$ as variables and insert this into the Einstein equations.

Setting $Bi = 0$, 4 of the 10 Einstein equations become trivial and one is left with 6 partial-differential equations (pdeq) in $r$ and $\theta$ for the 6 variables $A_0, \ldots, A_4$ and $M$.

The (outer) boundary condition here at the effective star radius $R$ with total mass $M_0$ is: $Ai = 1$, $M = M_0$, and $\partial_r Ai = 0$, $\partial_r M = 0$, as the density becomes 0 and the space-time becomes the normal Kerr space-time in vacuum.

Now, with rotation, we have a new model parameter, the angular velocity $\omega$, to which corresponds a third parameter in the outer boundary condition: (outer) ellipticity $\Delta R_1$, where $R_1 = R_\phi - \Delta R_1$ and $R_\phi$ and $R_1$ are the equatorial and the polar radius. As in the TOV-case, here to the 3 parameters $R_{ij}$, $M_0$ and $\Delta R_i$ correspond the 3 inner parameters $r_{ij}$, $\rho_0$ and $\Delta r_i$.

So here we get a 3-parametric solution manifold, and as in the spherical case, for a given total mass $M_0$ we have to find the stable physical solution. As before, these will be the ones with minimal $r_{ij}$ and among them the one with minimal mean energy density: this defines the inner ellipticity $\Delta r_i$. In all considered cases, it can be shown numerically, that such a (non-trivial) minimum exists.

The paper is organized as follows.

In 2 we present the mathematical setup, in 3 the equations for the extended Kerr space-time with rotation, in 4 the solution algorithm for it. In 5 the TOV-equation is introduced, in 6 the equation-of-state for the nucleon fluid and nucleon gas. In 7 the results for the TOV-equation are shown: the parametric solution manifold in 7.1. and the case study for typical stars in 7.2. In 8 the results for the extended Kerr space-time with rotation are presented for three typical star configurations: compact neutron star, stellar shell-star, and galactic shell-star.

2. The Kerr Space-Time, Schwarzschild Space-Time, Einstein Equations

Using the Minkowski metric $\eta_{\mu\nu} = \text{diag}(1, -1, -1, -1)$, the Kerr space-time metric in original Kerr coordinates $(u, \theta, \phi)$ has the line element \[^{[1]}\quad [3]\]

$$
\begin{align*}
\text{ds}^2 &= \left(1 - \frac{r_s}{r^2 + a^2 \cos^2 \theta}\right) \left[du + a \sin^2 \theta \, d\phi\right]^2 \\
&\quad - 2 \left(du + a \sin^2 \theta \, d\phi\right) \left(dr + a \sin^2 \theta \, d\phi\right) \\
&\quad - \left(r^2 + a^2 \cos^2 \theta\right) \left(d\theta^2 + \sin^2 \theta \, d\phi\right) \\
\end{align*}
$$

(1)

where $r_s = \frac{2GM}{c^2}$ is the Schwarzschild radius, and $a = \frac{J}{Mc}$ is the angular momentum radius (amr), $a$ has the dimension of a distance: $[a] = [r]$, and $J$ is the angular momentum.
With this line element the Kerr metric tensor $g_{\mu\nu}$ is as follows: [3]

$$g_{\mu\nu} = \begin{pmatrix}
1 - \frac{rr_s}{\rho_{12}} & -1 & 0 & \frac{-rr_s a \sin^2 \theta}{\rho_{12}} \\
0 & 0 & -a \sin^2 \theta & 0 \\
-\rho_{12} & 0 & -\rho_{12} \\
0 & 0 & -\rho_{12} & g_{33}
\end{pmatrix}$$

(2)

with the abbreviations $\rho_{12} = r^2 + a^2 \cos^2 \theta$ and $g_{33} = -\sin^2 \theta \left( r^2 + a^2 + \frac{rr_s a^2 \sin^2 \theta}{\rho_{12}} \right)$.

In the limit $a \rightarrow 0$ the Schwarzschild space-time in advanced Eddington-Finkelstein coordinates emerges:

$$ds^2 = \left(1 - \frac{r}{r_s} \right) du^2 - 2 du dr - r^2 \left( d\theta^2 + \sin^2 \theta d\varphi \right)$$

(3)

This form of the Schwarzschild line element has the advantage in comparison with the original line element

$$ds^2 = \left(1 - \frac{r}{r_s} \right) c^2 dt^2 - \frac{dr^2}{1 - \frac{r}{r_s}} - r^2 \left( d\theta^2 + \sin^2 \theta d\varphi \right)$$

(3a)

that the (apparent) singularity at $r = r_s$ is missing.

The same is valid for the original Kerr space-time: the denominator $\rho_{12}$ has no zeros, there is no singularity in $g_{ab}$, which makes it more well-behaved numerically.

Alternatively, in Boyer-Lindquist-coordinates: [3]

$$g_{\mu\nu} = \begin{pmatrix}
1 - \frac{rr_s}{r^2 + a^2 \cos^2 \theta} & 0 & 0 & \frac{rr_s a \sin^2 \theta}{r^2 + a^2 \cos^2 \theta} \\
0 & -\rho_{12}/\Lambda_{12} & 0 & 0 \\
-\rho_{12} & 0 & -\rho_{12} \\
0 & 0 & -\rho_{12} & -\sin^2 \theta \left( r^2 + a^2 + \frac{rr_s a^2 \sin^2 \theta}{\rho_{12}} \right)
\end{pmatrix}$$

(4)

with the line element

$$ds^2 = \left(1 - \frac{rr_s}{r^2 + a^2 \cos^2 \theta} \right) (dt)^2 + \left( \frac{2rr_s a \sin^2 \theta}{r^2 + a^2 \cos^2 \theta} \right) dt d\varphi - \left( \frac{r^2 + a^2 \cos^2 \theta}{r^2 - rr_s + a^2} \right) dr^2 - \left( r^2 + a^2 + \frac{rr_s a^2 \sin^2 \theta}{r^2 + a^2 \cos^2 \theta} \right) \sin^2 \theta d\varphi^2$$

(4a)

with the abbreviation $\Lambda_{12} = r^2 - rr_s + a^2$. Here, $\Lambda_{12}$ has zeros at the inner/outer horizon $r = (r_s/2) \pm \sqrt{(r_s/2)^2 - a^2 \cos^2 \theta}$, so for numerical calculations the singularity has to be removed by adding a small $\varepsilon$: $\Lambda_{12} = \sqrt{(r^2 - rr_s + a^2)^2 + \varepsilon^2}$.  

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In the limit $a \rightarrow 0$ the Schwarzschild space-time in the standard form (4) emerges.

The Einstein field equations with the above Minkowski metric are:

$$R_{\mu\nu} - \frac{1}{2} g_{\mu\nu} R - \Lambda g_{\mu\nu} = -\kappa T_{\mu\nu}$$

where $R_{\mu\nu}$ is the Ricci tensor, $R_0$ the Ricci curvature, $\kappa = \frac{8\pi G}{c^4}$, $T_{\mu\nu}$ is the energy-momentum tensor, $\Lambda$ is the cosmological constant (in the following neglected, i.e. set 0),

with the Christoffel symbols (second kind)

$$\Gamma_{\mu\nu} = \frac{1}{2} g^{\kappa\lambda} \left( \frac{\partial g_{\kappa\mu}}{\partial x^\nu} + \frac{\partial g_{\kappa\nu}}{\partial x^\mu} - \frac{\partial g_{\mu\nu}}{\partial x^\kappa} \right)$$

and the Ricci tensor

$$R_{\mu\nu} = \frac{\partial \Gamma^\rho_{\mu\nu}}{\partial x^\rho} - \frac{\partial \Gamma^\rho_{\mu\rho}}{\partial x^\nu} + \Gamma^\sigma_{\rho\nu} \Gamma^\rho_{\mu\sigma} - \Gamma^\sigma_{\rho\mu} \Gamma^\rho_{\sigma\nu}$$

The crucial part of the extended Kerr solution is the expression for the energy-momentum tensor $T_{\mu\nu}$. As usual, one uses the formula for the perfect fluid [1, (45.3)]:

$$T_{\mu\nu} = \left( \rho + \frac{P}{c^2} \right) u_\mu u_\nu - P g_{\mu\nu}$$

where $P$ and $\rho$ is the pressure and density, $u_\mu$ is the covariant velocity 4-vector.

In the Schwarzschild case, when deriving the TOV-equation, one sets the spatial contravariant velocity components to 0: $u^t = 0$ , in the Kerr case the tangential velocity $u^3 = u^\phi \neq 0$.

For the velocity one has:

$$u^3 = r \omega = r \frac{aMC}{I}$$

where $I$ is the moment of inertia and $\omega$ the angular velocity,

$$I = f_1 MR_i^2 , \text{ so}$$

$$u^3 = \frac{1}{2} \frac{acr}{f_1 R_i^2}$$

for a homogeneous sphere $I = \frac{2}{5} MR^2$ i.e. $f_1 = \frac{2}{5}$, for a thin shell $I = \frac{2}{3} MR^2$ i.e. $f_1 = \frac{2}{3}$, for a disc $I = \frac{1}{2} MR^2$.

If we make the obvious assumption that the star rotates as a whole, i.e. with constant angular velocity, then the moment of inertia $I$ becomes $r$-dependent, like the mass $M$:

$$M(r) = \int_0^r \rho(r) 3r^2 dr$$

$$I(r) = \int_0^r \rho(r) 3r^4 dr$$

The factor 3 in the integral instead of the usual $4\pi$ comes from the dimensionless calculation in “sun units” (see below).
The parameter \( a \) also becomes \( r \)-dependent:

\[
a(r) = \frac{J}{Mc} = \frac{I(r)\omega}{M(r)c}
\]

In the relativistic axisymmetric case with rotation with angular velocity \( \omega \), \( u^\nu \) has the form [4]:

\[
u = u^0 (1,0,0,\omega)
\]

\[
u = \left(u_0,0,0,\omega u_0\right)
\]

Now \( u^0 \) is calculated from the condition

\[
c^2 = g_{\mu\nu}u^\mu u^\nu
\]

and the covariant velocity from

\[
u_\mu = g_{\mu\nu}u^\nu
\]

The resulting expression for \( u^0 \) is (\( A_i \) are the Kerr correction-factors, mass \( M_1[r_1] \), moment of inertia \( I_1[r_1] \)):

\[
u^0 = \left(A_i(r_1,\theta)\left[\frac{r_1M_1(r_1)^3}{\omega^2I_i(r_1)^2\cos^2(\theta) + r_1^2M_1(r_1)^2} - 1\right]ight)
\]

\[
- \frac{r_1\omega^2I_i(r_1)^2M_1(r_1)\cos^4(\theta)A_i(r_1,\theta)}{\omega^2I_i(r_1)^2\cos^2(\theta) + r_1^2M_1(r_1)^2}
\]

\[
- \frac{\omega^4I_i(r_1)^2\sin^2(\theta)A_i(r_1,\theta)}{M_1(r_1)^2}\frac{r_1^2\omega^2\sin^2(\theta)A_i(r_1,\theta)}{M_1(r_1)^2}
\]

\[
+ \frac{2\omega^2\sqrt{I_i(r_1)^2M_1(r_1)^2\sin^2(\theta)A_i(r_1,\theta)}}{\omega^2I_i(r_1)^2\cos^2(\theta) + r_1^2M_1(r_1)^2}^{-1}
\]

The state equation for the pressure \( P \) for the nucleon gas has the form

\[
P = c_1\rho^\gamma
\]

or in the dimensionless form with a critical density \( \rho_c \) and dimensionless pressure \( P_1 \) and density \( \rho_1 \)

\[
P_1 := \frac{P}{c^2\rho_c^\gamma} = k_1\left(\frac{\rho}{\rho_1}\right)^\gamma = k_1(\rho_1)^\gamma
\]

For the horizon, with rotation there is the inner and the outer horizon \( (M = M_0) \)

\[
r_+ = M_0 \sqrt{\left(\frac{M_0}{2}\right)^2 - a^2}
\]

\[
r_+ = M_0 \sqrt{\left(\frac{M_0}{2}\right)^2 + a^2}
\]

3. The Equations for the Extended Kerr Space-Time

The solution process starts with the metric tensor \( g_{\mu\nu} \) in original Edding-
ton-Finkelstein-coordinates, with 6 non-zero components, corresponding correction-factor functions $A_0, \ldots, A_5$, and additive correction functions $B_0, \ldots, B_3$ for the zero components.

$$
g_{\mu\nu} = \begin{pmatrix} 
1 - \frac{r r}{\rho_{12}} & A_0 & -A_3 & B_1 & -\frac{r r a \sin^2 \theta}{\rho_{12}} & A_4 \\
B_0 & B_2 & -A_2 \rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_3 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
-A_2 \rho_{12} & 1 - \frac{r r a \sin^2 \theta}{\rho_{12}} & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 \\
-A_2 \rho_{12} & 1 - \frac{r r a \sin^2 \theta}{\rho_{12}} & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 \\
-A_2 \rho_{12} & 1 - \frac{r r a \sin^2 \theta}{\rho_{12}} & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 \\
-A_2 \rho_{12} & 1 - \frac{r r a \sin^2 \theta}{\rho_{12}} & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 \\
\end{pmatrix}$$

and alternatively in Boyer-Lindquist-coordinates, corresponding correction-factor functions $A_0, \ldots, A_4$, and additive correction functions $B_0, \ldots, B_4$ for the zero components

$$
g_{\mu\nu} = \begin{pmatrix} 
1 - \frac{r r}{\rho_{12}} & A_0 & B_0 & B_1 & -\frac{r r a \sin^2 \theta}{\rho_{12}} & A_4 \\
-\rho_{12} A_1 & B_2 & -\rho_{12} A_2 & -\rho_{12} & -\rho_{12} A_1 & B_3 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
\end{pmatrix}$$

The equations are the 10 Einstein equations eqR00, eqR11, eqR22, eqR33, eqR12, eqR23, eqR31, eqR01, eqR02, eqR03 in the (dimensionless) variables relative radius $r_1 = \frac{r}{r_{ss}}$ and complementary azimuth angle $\theta_1 = \frac{\pi}{2} - \theta$ with energy tensor $T_{\mu\nu}$ from (8) and the state equation $P_1 = k_1 (\rho_1)^\gamma$ for the relative pressure $P_1$ and the relative density $\rho_1$. We are using the so called “sun units” $r_{ss} = r_1 (\text{sun})$, $M_s = M (\text{sun})$, $\rho_s = \frac{M}{4\pi r_{ss}^3}$, $P_s = \rho_s c^2$ for radius $r$, mass $M$, density $\rho$ and pressure $P$, respectively.

In “sun units” the original angle differential $d\Omega = 4\pi \sin \theta r^2 d\theta d\phi$ is transformed into $d\Omega = 3 \cos \theta r^2 d\theta d\phi$, as for $\theta = 0, \ldots, \pi/2$, $r = 0, \ldots, 1$: $\int d\Omega = 1$.

Also, all equations and variables are symmetric (even) in $\theta$. $A(\theta) = A(\pi - \theta)$.

From now on we skip the index of the dimensionless variables and use the original notation, e.g. $r$ instead of $r_1$.

Furthermore, we adopt the Boyer-Lindquist coordinates and the metric tensor (12).

In “sun units”, the Boyer-Lindquist metric tensor becomes:

$$
g_{\mu\nu} = \begin{pmatrix} 
1 - \frac{r M_0}{\rho_{12}} & A_0 & B_0 & B_1 & -\frac{r M_0 a \cos^2 \theta}{\rho_{12}} & A_4 \\
-\rho_{12} A_1 & B_2 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
-\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & -\rho_{12} & -\rho_{12} A_1 & B_4 \\
\end{pmatrix}$$
\[
\rho_{12} = r^2 + a^2 \sin^2 \theta \\
\Lambda_{12} = r^2 - rM_0 + a^2
\]

where \( M_0 \) is the mass in sun units.

The 10 Einstein equations have a distinctive structure: there are 6 primary variables \( A_0, A_2, A_3, A_4, B_1, B_4 \) with highest derivative \( \partial_r \) and 4 secondary variables \( A_1, B_2, B_0, B_3 \) with highest derivative \( \partial_\theta \). Primary variables have boundary conditions for the variable and its \( r \)-derivative, secondary variables only for the variable itself.

This structure is dual in \( \theta \): again there are 6 \( \theta \)-primary and 4 \( \theta \)-secondary variables.

6 of Einstein equations contain only one 2-derivative \( \partial_r \) of a primary variable: \( \text{eqR03}(\partial_r A_4) \), \( \text{eqR22}(\partial_r A_2) \), \( \text{eqR00}(\partial_r A_0) \), \( \text{eqR33}(\partial_r A_3) \), \( \text{eqR02}(\partial_r B_1) \), \( \text{eqR23}(\partial_r B_4) \) contain only 2-derivatives of a secondary variable: \( \text{eqR12}, \text{eqR01}, \text{eqR31}, \text{eqR11} \) contains all derivatives \( \partial_r \) of a primary variable.

If we make the ansatz \( B_i = 0 \), several of the eqRij become identically 0, and we get the 6 equations eqR00, eqR11, eqR22, eqR33, eqR03, eqR12 for the 6 variables \( A_i \) and \( \rho \), with the highest derivatives resp. \( \partial_r A_0, \partial_{\theta\theta} A_1, \partial_r A_2, \partial_r A_3, \partial_r A_4, (\partial_\theta A_2, \partial_{\theta\theta} A_1) \).

Thus, we are left with the 6 differential equations degree 2 in \( r, \theta \) non-linear (quartic) in variables \( A_i \) and their 1-derivatives and linear in \( \rho, \rho^\gamma \).

In total, we have 6 algebro-differential eqs for 6 variables \( A_i \) and \( \rho \) (\( \rho \) enters only algebraically).

We can add 2 dependent equations eqR41 \( \equiv D_\mu T^{\mu i} = 0 \) and eqR42 \( \equiv D_\mu T^{\mu 2} = 0 \), from the covariant continuity equations \( D_\mu T^{\mu i} = 0 \), where \( D_\mu T^{\mu i} = \partial_\mu T^{\mu i} + \Gamma^\kappa_\mu_\lambda T^{\kappa i} + \Gamma^\kappa_\lambda_\mu T^{\mu i} \) is the gravitational covariant derivative.

In eqR41 \( \rho \) enters with \( \partial_r \rho \), in eqR42 \( \rho \) enters with \( \partial_\theta \rho \).

So, alternatively, we have the diff. equations eqR00, eqR11, eqR22, eqR33, eqR03, eqR41, with the highest derivatives resp. \( \partial_r A_0, \partial_{\theta\theta} A_1, \partial_r A_2, \partial_r A_3, \partial_r A_4, \partial_\theta A_2, \partial_{\theta\theta} A_1 \) (diff. eq. degree 1 in \( r \) for \( \rho \)) or the diff. equations eqR00, eqR11, eqR42, eqR33, eqR03, eqR42, with the highest derivatives resp. \( \partial_r A_0, \partial_{\theta\theta} A_1, \partial_r A_2, \partial_r A_3, \partial_r A_4, \partial_\theta A_2, \partial_{\theta\theta} A_1 \) (diff. eq. degree 1 in \( \theta \) for \( \rho \)).

In the Schwarzschild spacetime \( a = 0 \) and \( a = 0 \), we have spherical symmetry, no dependence on \( \theta \), and the TOV-equation can be derived from the non-trivial eqR00, eqR11, eqR42, eqR41.

We impose an \( r, \theta \)-analytic boundary condition for \( A_i, \partial_r A_i \) at \( r = R_1 \) (\( R_1 \) is the star radius):

\( A_i = 1, \partial_r A_0 = 0, \partial_r A_2 = 0, \partial_r A_3 = 0, \partial_r A_4 = 0 \). For \( A_1 \), there is no differential boundary condition, as \( \partial_r A_1 \) is the highest \( r \)-derivative, for \( \rho \) there is no boundary condition at all, because \( r \) is algebraic in the equations, but there is an integral condition:

\[
M(R_i) = \int_0^{\sqrt{\sigma}} \rho(r, \theta) 3r^2 \cos \theta \, dr \, d\theta = M_0 : \text{integral(} \rho \text{) = total mass} = M_0
\]
In order to avoid the clumsy integral condition for $\rho$, we can introduce the mass $M$ as a variable:

$$M(r, \theta) = \int_0^r \rho(\bar{r}, \theta) 3\bar{r}^2 d\bar{r}$$

$$M(r) = \int_0^{\pi/2} M(r, \theta) d\theta = \int_0^r \int_0^{\pi/2} \rho(\bar{r}, \theta) 3\bar{r}^2 d\bar{r} \cos \theta d\theta$$

is the mass of the sphere($r$) and

$$\partial_r M(r, \theta) = \rho(r, \theta) 3r^2$$

For $M(r, \theta)$ we impose the boundary condition at $r = R_1$:

$$M(R_1, \theta) = M_0, \quad \partial_r M(R_1, \theta) = 0 \quad (i.e. \ density \ \rho \ is \ zero \ at \ boundary, \ and \ total \ mass \ M_0).$$

So, if we take the diff. equations eqR00, eqR11, eqR22, eqR33, eqR41 and replace $\rho(r, \theta)$ by $\partial_r M(r, \theta)$, we have 6 diff.equations in $r, \theta$ of degree 2, for the variables $A_i(r, \theta)$ (metric correction factors = mcf) and the mass $M(r, \theta)$, with the highest derivative $\partial_{rr} M(r, \theta)$ in $M$.

According to the Cauchy-Kovalevskaya theorem there exists then a unique solution in a region $R_i r > 0$ within the boundary. Inside the region $r_i > r > 0$ we can enforce the vacuum Kerr-spacetime with the trivial solution $A_i = 1, \rho = 0, i.e. \ there \ is \ no \ matter \ there, \ r_i \ the \ inner \ radius.$

The Cauchy-Kovalevskaya theorem guarantees the existence of a mathematical solution outside the horizon, but for a physical solution we must have $r \geq 0$ (meaning $\partial_r M(r, \theta) > 0$) and $M(r, \theta) \leq 0$ for $r \leq r_i$: the mass must become non-positive at the inner radius.

Therefore, for certain {$M_o, R_1$} values there will be no physical solution, even for the TOV equation.

4. The Solving Process for the Extended Kerr Space-Time

In addition to the fundamental dual parameters {$r_i, \rho_i$} corresponding to {$R_1, M_0$} in the rotation-free TOV-case, in the Kerr-case there is the new fundamental parameter $\Delta r_i$ (inner ellipticity for inner boundary condition), resp. $\Delta R_1$ (outer ellipticity for outer boundary condition), and the angular velocity $\omega$. The outer radii are

$$R_{11} = R_i - \Delta R_i \quad \text{and} \quad R_{s1} = R_1,$$

the latter equality arising from the fact that centrifugal distortion acts only in the x-direction (the y-axis being the rotation axis). The inner radii are correspondingly $r_{s1} = r_i - \Delta r_i$, $r_i = r_{s1}$.

The $r-\theta$-slicing algorithm with an Euler-step obeys the iterative procedure with slice step size $h_r$ in $r$, and step size $h_\theta$ in $\theta$, starting with the $r$-boundary at $r = R_1$ (slice $n = 0$).

The transition from slice $n$ to $n + 1$ proceeds as follows.

At slice $n$ all variables and 1-derivatives are known from the previous step, 2-derivatives $\partial_{rr} A_i$, and $\rho$ are calculated from the 6 equations.

At slice $n + 1$ the variables and 1-derivatives are calculated by Euler-formula (or Runge-Kutta)

$$A_i_{n+1} = A_i_n + h_\theta \partial_r A_i_n$$
\[ \partial_r A_i = \partial_r A_i + h_0 \partial_r A_i \]

The 2-derivatives \( \partial_r A_i \), \( \partial_r B_i \) and \( \rho \) are again calculated from the 6 significant equations with variables and 1-derivatives inserted from above.

The \( \theta \)-slicing-backward algorithm with an Euler-step obeys the iterative procedure with slice step size \( h_0 \) in \( \theta \) as above for \( r \), starting with \( \theta = 0 \), and solves an ordinary differential equation in \( r \) in each \( \theta \)-step. The boundary condition for the \( r \)-odeq is set at \( r = R_1(\theta) \) (the outer ellipse radius) with \( A_i = 1 \), \( M = M_0 M(0)(\theta) \), \( \partial_r A_i = 0 \), \( \partial_r M = 3(R_1^2) \rho_{bc} \), where \( \rho_{bc} \) is the outer boundary value for the density, \( \rho_{bc} = 0 \) for the (non-interacting) neutron-gas in a shell-star and \( \rho_{bc} > 0 \), \( \rho_{bc} = \rho_{\text{equilibrium}} \) for the (interacting) neutron fluid in a neutron star. \( M_0(\theta) \) is the mass-form-factor with the condition \( \int_0^{2\pi} M_0(\theta) \cos(\theta) d\theta = 1 \), i.e. the overall mass at the outer boundary is \( M_0 \). The inner radius \( r_i(\theta) \) is reached, when \( M_0(\theta) = 0 \).

The alternative (dual) \( \theta \)-slicing-forward algorithm starts with the boundary condition at

\[ r = r_i(\theta) \]

\( A_i = 1 \), \( M = 0 \), \( \partial_r A_i = 0 \), \( \partial_r M = 3(r_i^2) \rho_{bc} \), where \( \rho_{bc} = \rho_i \) is the inner boundary value for the density, \( \rho_i \) is approximately the inner (maximum) density \( \rho(r_i) \) from the corresponding TOV-equation, the value must be adapted, so that the resulting total mass is \( M_0 \). For the compact neutron star the inner radius \( r_i(\theta) \) is zero.

In the \( \theta \)-slicing-backward algorithm one starts with the outer boundary being the ellipsoid \( r = R(\theta, \Delta R_1) \), where \( \Delta R_1 \) is the outer ellipticity of the star. In the \( \theta \)-slicing-forward algorithm one starts with the inner boundary being the ellipsoid \( r = r_i(\theta, \Delta r_i) \), where \( \Delta r_i \) is the inner ellipticity of the star.

At the inner boundary the tangential pressure is uniform, so the density is also uniform and equal to the maximum density, \( \rho(\theta) = \rho_c \).

In the actual calculation we were using the \( \theta \)-slicing-backward algorithm, because here the boundary condition \( M = M_0 \) is achieved automatically, when one starts with \( M_0(\theta) = M_0 \).

The odeqs in \( r \) consist of the 6 significant Einstein equations \( \text{eqR00, eqR11, eqR22, eqR33, eqR03, eqR41} \) for the six variables \( A_0(r, \theta), A_1(r, \theta), A_2(r, \theta), A_3(r, \theta), A_4(r, \theta), M(r, \theta) \) with \( \theta = \theta \) and \( \theta \)-derivatives calculated by Euler-step from the preceding q-slice. For \( i = 0 \) i.e. \( \theta = 0 \) the \( \theta \)-derivatives are taken from start values for all variables, which normally represent the corresponding TOV-solution (here only \( A_0(r), A_1(r), M(r) \) are non-trivial and do not depend on \( \theta \)). The odeqs are highly non-linear algebraic differential equations and hard to solve numerically with classical methods for linear odeqs extended by an algebraic equation solver. In the case of a nonlinear odeq-system one uses an Euler or Runge-Kutta method and calculates in each step the highest derivatives with a numerical algebraic equation solver. As an alternative one can use minimization of the least-squares-error in the highest derivatives instead of a numerical algebraic equation solver. Minimization has also the advantage that one can minim-
ize the complete set of Einstein equations plus the 2 additional continuity equations eqR41, eqR42 in the error goal function instead of the 6 significant equations, which improves the stability of the solution (e.g. in case of degeneracy).

The numerical error of the algorithm is calculated from \( \sum_i (eq_i)^2 \), i.e. the Euclidean norm of the equation values (the right side of the Einstein equations being 0). The error is calculated over the lattice \( \{r_i, \theta_j\} \) as median, mean or maximum. In the internal loop of the algorithm over \( r \), at fixed \( \theta \), the solution of the algebraic discretised Einstein-eqation is achieved by square-root error minimization, so it is essential to avoid singularities, e.g. at the horizon and the pseudo-singularity at \( \theta = 0 \). This is achieved by selecting appropriate analytic convergence factors for the (left side of) the Einstein equations. As the equations are to be zeroed for the solution, the convergence factors do not change the solution of course, but they cancel the numerical singularities, which could otherwise jeopardize the numerical convergence of the algorithm.

The actual calculation was carried out in Mathematica using its symbolic and numerical procedures. In the first stage, the Einstein equations were derived from the ansatz for \( g_{\mu\nu} \) from section 2 and simplified automatically. The arising complexity of the equations is such, that it is practically impossible to handle them manually: the Mathematica function LeafCount, which returns the number of terms in the equation, gives the complexity of \( \text{LeafCount} \text{[eqR00]} = 17,408 \), \( \text{LeafCount} \text{[eqR11]} = 27,528 \), \( \text{LeafCount} \text{[eqR22]} = 134,929 \) for the first 3 equations. To verify the equations, the TOV equation was derived by symbolic manipulation for \( \omega = 0 \ a = 0 \) from eqR00, eqR11, eqR22, eqR41.

The power of Mathematica is sufficient to solve the TOV equation with the single procedure NDSolve. For the full Einstein equations it fails even for the ordinary differential equations (odeq) in arising for fixed \( \theta \). It took us a long time to find an algorithm, which could handle the complexity of the equations and solve them in an acceptable time (3 - 5 minutes on a 16 \times 8 lattice) on a PC-desktop and converge in the required region with an acceptable error of around 0.01.

For the second numerical stage we tried several slicing algorithms, and the best alternative proved to be the \( \theta \)-slicing-forward algorithm implemented by hand in Mathematica. The solution of the resulting odeq in each r-step was calculated using NDSolve.

Also, for every star model and parameter set, the TOV solution with \( \omega = 0 \ a = 0 \) was calculated first with the algorithm and compared with the exact TOV solution.

5. The TOV Equation as the Limit \( \omega \to 0 \) for the Extended Kerr Space-Time

In the Schwarzschild spacetime \( \omega = 0 \) and \( a = 0 \), we have spherical symmetry, no dependence on \( \theta \), then the TOV-equation can be derived from the remaining non-trivial Einstein equations eqR00, eqR11, eqR22, eqR41.

The TOV-equation is in the standard form:
\[ P'(r) = -\left( GM(r) \rho(r) \right) \left( \frac{1 + P(r)}{\rho(r)c^2} \right) \left( 1 - \frac{2GM(r)}{rc^2} \right)^{-1} \]

and using \( r_s \)

\[ P'(r) = -\left( \frac{c^2r_s\rho(r)}{2r^2} \right) \left( \frac{1 + P(r)}{\rho(r)c^2} \right) \left( 1 - \frac{r_sM(r)}{rM_i} \right)^{-1}, \]

where

\( M_i \) is the total mass, furthermore

\[ 4\pi r^2 \rho(r) = M(r), \quad P(r) = k_i \rho(r)^\gamma \]

In order to make the variables dimensionless, one introduces “sun units”

\[ r_{ss} = r_s (\text{sun}) = \frac{2GM_{\odot}}{c^2} = 3 \text{ km}, \quad \rho_s = \frac{M_{\odot}}{4\pi r_{ss}^3/3} = 1.76 \times 10^{16} \frac{g}{\text{cm}^3}, \quad P_s = \rho_s c^2 \]

where \( r_{ss} \) Schwarzschild-radius of the sun, \( \rho_s \) the corresponding Schwarzschild-density and \( P_s \) the corresponding Schwarzschild-pressure.

In “sun units” TOV-equation transforms into

\[ P_1'(r_1) r_1^3 \left( r_1 - M_i(r_1) r_0 \right) = -\frac{1}{2} \left( \frac{M_1'(r_1) r_0}{3} + P_1(r_1) r_1^3 \right) (M_1(r_1) r_0 + 3P_1(r_1) r_1^3) \]

with the normalized mass \( M_i(r_i) \), and \( M_i(R_i) = 1 \), or

\[ P_1'(r_1) r_1^3 \left( r_1 - M_i(r_1) \right) = -\frac{1}{2} \left( \frac{M'(r_1)}{3} + P_1(r_1) r_1^3 \right) \left( M(r_1) + 3P_1(r_1) r_1^3 \right) \]

where \( M_0 = \frac{M_i}{M_{\odot}}, M_i(r_i) \) the mass within the radius \( r \), \( M_i(r_i) = M_iM_i(r_i) \)

in dimensionless variables \( r_i \), \( \rho_i(r_i) = \frac{M'(r_i)}{3r_i^2} \), \( M, \quad P_i = k_i \rho_i(r_i)^\gamma \) and \( R_i \) is the dimensionless radius of the star.

With the replacement \( P = k_i \rho^\gamma \) for the pressure from the equation of state and \( \rho = \frac{M_0M'}{3r^2} \), we obtain a diff. equation for \( M \) degree 2 in \( r \) and we impose the boundary condition in \( r = R_i \):

\[ M(R_i) = M_0, \quad M'(R_i) = 0 \]

for non-interacting Fermi-gas and for an interacting Fermi-gas:

\[ M(R_i) = M_0, \quad M'(R_i) = \rho_i(R_i) 3R_i^2, \quad \rho_i(R_i) = \rho_e, \]

where \( \rho_e \) is the equilibrium density in the minimum of \( V_{ss} \) and \( P_i'(\rho_e) = 0 \) (here an equivalent boundary condition is \( \rho'(R_i) = \infty \)).

6. The Equation of State and Rotation Parameters

6.1. The Equation of State for an (Non-Interacting) Nucleon Gas

Here, \( P = k_i \rho^\gamma \) is the equation of state of the star, derived from the thermodynamic Fermi gas equation at \( T = 0 \) ([1], chap. 48).
\[
P = -\frac{\partial E}{\partial V} = 8\pi P_0 \left( \frac{x_F^2}{3} \sqrt{1 + x_F^2} - f(x_F) \right)
\]  

(15)

\[
P_0 = \frac{mc^2}{\lambda_c^3} = \frac{m^4 c^4}{\hbar^3}, \text{ where } \lambda_c \text{ is the de-Broglie wavelength of the Fermi gas with particle mass } m, \quad \lambda_c = \frac{\hbar}{mc}.
\]

\[
x_F = \frac{P_F}{mc} = \frac{\lambda_c}{2} (3\pi)^{1/3} n^{1/3}, \text{ where } x_F \text{ is the Fermi-angular-momentum, } n \text{ the particle density}
\]

\[
f(x_F) = \int_0^{x_F} dx x^2 \sqrt{1 + x^2}
\]

The resulting approximate equations of state for \( P \) are

\[
P = 8\pi P_0 \begin{cases}
\frac{x_F^2}{15} & \rho \ll \rho_c \\
\frac{x_F^2}{12} & \rho \gg \rho_c
\end{cases}
\]

\[\begin{cases}
K_1 \rho^{2/3} & \rho \ll \rho_c \\
K_2 \rho^{2/3} & \rho \gg \rho_c
\end{cases}
\]

valid for the density \( \rho \) and the critical density \( \rho_c \)

\[
\rho_c = \frac{m 8\pi}{\lambda_c^3} 3 
\]

The full expression for \( P \), including temperature \( T \), is as follows ([5], chap.15).

Here, we use dimensionless variables \( r_1 \text{ distance unit de-Broglie-wavelength } \lambda, V_1 \text{ volume unit } \lambda_c^3, n_1 \text{ particle density unit } 1/\lambda_c^3, E_1 \text{ energy unit } E_0 = \frac{\hbar c}{\lambda_c^2} = \frac{mc^2}{2\pi}, \)

inverse thermal energy \( \beta_1 = \frac{E_0}{kT}, \) chem. potential \( \mu_1 \text{ in } E_0 \), for the gas model we use the Debye model with the state density

\[
\epsilon_{F_1} = \frac{3^{2/3} \pi^{2/3}}{4} n_1^{2/3}, \text{ the resulting particle density is}
\]

\[
n_1 = \frac{N_{\text{phys}}}{V_1} = \frac{2}{V_1} \int_0^{E_1} d\omega_1 \frac{D_1(\omega_1)}{1 + \exp(\beta_1(\omega_1 - \mu_1))}
\]

\[
= \frac{1}{2\pi^{1/2}} \int_0^{\pi/2} d\omega_1 \sqrt{\omega_1} \frac{1}{1 + \exp(\beta_1(\omega_1 - \mu_1))}
\]

From this relation the chem. potential \( \mu_1 \) can be calculated, an approximation formula is

\[
\mu_1 = \epsilon_{F_1} - \frac{\pi^2}{12 \beta_1^2 \epsilon_{F_1}} = \mu_1(n_1)
\]

Finally, the resulting pressure (=energy density) \( p_1(\beta_1, n_1) \):

\[
p_1(\beta_1, n_1) = \frac{4\pi^{2/3}}{3\pi^{1/2}} \int_0^{\pi/2} d\omega_1 \frac{\omega^{2/3}_1}{1 + \exp(\beta_1(\omega_1 - \mu_1(n_1)))}
\]

(17)

Below a 3D-diagram of \( p_1(\beta_1, n_1) \) in dimensionless variables for a nucleon
gas \( m = m_n \) density \( \rho = \frac{E_0 n}{c^2} \) in sun units, \( E_0 = 149.4 \text{ MeV} \) is depicted in

**Figure 1** [6].

Here \( kT \) is in \( E_0 \) units, and one sees the dependence \( P = kT \rho \) except on the left side, when \( kT \) reaches the magnitude of 1 Gev \((T = 10^{10} \text{K})\).

### 6.2. The Equation of State for an (Interacting) Nucleon Fluid

For the interacting nucleon gas we take into account the nucleon-nucleon-potential in the form of a Saxon-Woods-potential modeled on the experimental data: [7]-[13]

\[
V_{nn}(r, r_0, d_0) = \frac{V_0}{1 + \exp\left(\frac{r - r_0}{d_0}\right)}
\]

\[
P_{nn}(r) = -V_{nn}(r, r_0, d_0) + V_{nn}(r_0, r, d_0)
\]

where \( V_{nn} \) is the nucleon-nucleon-potential with an attractive part

\[
V_{nn}(r, r_0, d_0) = -V_{nn}(r_0, r, d_0)
\]

and a repulsive core \( V_{nn}(r_0, r, d_0) \), the distance \( r \) between the nucleons is \( r = (E_n/\rho)^{1/3} \), where \( E_n = (m_n c^2 / (2\pi))^{1/3} = 149.4 \text{ MeV} \approx m_n c^2 \)

is the nuclear energy scale \( m_n = \) pion mass = 140 MeV, \( m_n = \) neutron mass = 140 MeV.

The Saxon-Wood potential is shown in **Figure 2** below.

The pressure of the interacting nucleon fluid becomes then

\[
P_{nn}(r) = e(r) \rho_1(r) V_{nn}\left((E_n/\rho_1(r))^{1/3}\right)
\]

The experimental data used here are those from [7], and are shown in **Figure 3**.

And the hard-core potential from the lattice calculation Reid93 [10] is shown in **Figure 4**.

Both potentials are fitted with a double Saxon-Woods-potential \( V_{nn} \) in **Figure 5**.

**Figure 1.** Fermi pressure.
From the nucleon-nucleon-potential the pressure is calculated taking into account the low-density Fermi-pressure of the nucleons \( P_{\text{ff}}(\rho) = K_1 \rho^{5/3} \) and nucleon-nucleon pressure \( P_{\text{nn}}(\rho) = V_{\text{nn}}(\rho^{1/3}) \rho \) is shown in Figure 6 [6], and the total pressure \( P_{\text{gg}}(\rho) = K_1 \rho^{5/3} + P_{\text{nn}}(\rho) \) is shown in Figure 7.

This equation-of-state has a minimum at \( \rho = \rho_c = 0.0417 \) and \( P'(r) = 1 \) at \( \rho = \rho_m = 0.0544 \).
Figure 5. Nucleon-nucleon potential data fitted with Saxon-Wood-formula, with \( r(\text{fm}) \), \( V(\text{MeV}) \).

Figure 6. Nucleon-nucleon Fermi-pressure.

Figure 7. Total pressure \( P_g(r) \), pressure \( P \) and density \( \rho \) shown in sun-units.
As the sound velocity \( v = \frac{dP(\rho)}{d\rho} \), \( v > 0 \) and \( v < 1 \) (i.e. subluminal), the admissible density range in the neutron-fluid model is \( \rho_c \leq \rho \leq \rho_w \).

### 6.3. Maximum Omega-Values in Kerr-Space-Time

We consider here a rotation model with constant angular velocity \( \omega \). With this model the resulting 4-velocity \( u^\mu \) has the form \([4][14][15]\):

\[
\begin{align*}
  u^\mu &= (u^0, 0, 0, \omega u^0) \\

\end{align*}
\]

The maximum values for \( \omega \) are calculated from the minimal zeros in omega of the denominator in \( u^0 \) from (9a), minimized over \( r_1 \) and \( \theta \) in their respective regions

\[
  r_i \leq r \leq R_i \quad \text{and} \quad 0 \leq \theta \leq \pi/2.
\]

The resulting value is

\[
  \omega \leq \frac{1}{2R_i \sqrt{\alpha_f}},
\]

where \( \alpha_f \) is the form-factor in the moment of inertia \( I_1 \).

\[
  I_1 = \alpha_f M R_1^2, \quad \alpha_f = 2/3 \text{ for a shell, } \alpha_f = 2/5 \text{ for a sphere.}
\]

With non-vanishing density the actual \( \omega_{\text{max}} \) depends on \( \{A, \rho\} \), and has to be calculated from the above expression for \( u^0 \).

A less stringent limit for omega can be deduced from the limit for the parameter \( a \) in Kerr-spacetime (in sun-units and \( c = 1 \)): \( a \leq \frac{r_1 (M_0)}{2} = \frac{M_0}{2} \), and

\[
  a = \frac{\alpha_f M R_1^2 \omega}{M_0} \quad \text{therefore} \quad \omega \leq \frac{M_0}{2R_1^2 \alpha_f}.
\]

### 7. The TOV-Equation: A New Ansatz

Generally speaking, the parameters of the solution are (in brackets denomination in program code): angular momentum radius \( a (= \text{alpha1}, \text{=0 for TOV}) \), the factor in the state equation \( k_1 \), the power in the state equation \( \gamma (= \text{gam}) \), radius \( R \), mass \( M_0 \), the relative radius uncertainty \( dr_{\text{rel}} \), \( (= \text{dr02rel}) \), the moment of inertia factor \( f_i (= \text{infac}, \text{=0 for TOV}) \), the singularity smoothing parameter \( \varepsilon (= \text{epsi}, \text{see below}) \), and the boundary factor \( n_{\text{max}} (= \text{nrmax}) \). Here the boundary factor enters the upper boundary of the TOV differential equation as \( r_{\text{max}} = R(1 + n_{\text{max}} dr_{\text{rel}}) \).

The dimensionless TOV-equation is a differential equation in the mass \( M(r) \) of degree 2, and is highly non-linear, the dimensionless mass-density relation is

\[
  \rho = \frac{M}{3r^2}.
\]

The customary way of solving the TOV equation is to impose the boundary condition at \( r = 0 \) with \( M(0) = 0 \), \( M'(0) = 3r^2 \rho_0 \), where \( \rho_0 \) the maximum central density.

In the new ansatz for the mass \( M(r) \) we impose the outer boundary condition at \( r = R_i \):
for a pure Fermi-gas without interaction: \( M(R_0) = M_0, \ M'(R_0) = \rho(R_0) 3R_0^2, \rho(R_0) = 0; \)

for an interacting Fermi-gas: \( M(R_0) = M_0, \ M'(R_0) = \rho(R_0) 3R_0^2, \rho(R_0) = \rho_e, \) where \( \rho_e \) is the equilibrium density in the minimum of \( V_{\text{nn}} \) and \( P'\left(\rho_e\right) = 0 \) (here an equivalent boundary condition is \( P'(R_1) = \infty \).

The star parameters mass \( M_0 \) and radius \( R_1 \), which enter the outer boundary condition determine completely the solution. In general, there will be an inner radius \( r_i > 0 \) with the maximum density \( \rho_0 = 3r_i^2M'(r_i) \) and \( M(r_i) = 0 \). The corresponding “dual” parameters are the inner radius \( r_i \) and the maximum density \( \rho_0 \). One can show that for \( \rho_0 \gg \rho_c \) (where \( \rho_c \) is the critical density of the equation of state) there is no solution with a compact star \( r_i = 0 \), i.e. there is a maximum mass \( M_c \) for the TOV equation, in case of compact neutron stars \( M_c = 3.04\, M_{\odot} \) (see below). As we will see, there is in general a solution, if we allow \( r_i > 0 \) and impose an outer boundary condition at \( r = R_1 \), as long as \( R_1 \) is not too close to the Schwarzschild radius \( r_s = M_0 \) of the star. In the limit \( R_1 \to r_i \) there will be no positive zero of \( M(r) \), i.e. \( r_i < 0 \) and the resulting (mathematical) TOV-solution will be no physical solution. But in general, speaking naively, the gravitational collapse of the star is avoided for large masses \( M_0 > M_c \), if it has a shell structure with the inner radius \( r_i \) and the outer radius \( R_1 > M_0 \).

As we will see, this outer boundary condition together with allowing \( r_i > 0 \) changes dramatically the resulting manifold of physical solutions.

7.1. The TOV-Equation: The Parametric Solution and Resulting Star Types

By setting-up a parametric solution of the TOV-equation one gets a map of possible physical solutions, i.e. possible star structures. As parameters one can use either \( (M_0, R_1) \) in the outer boundary condition at \( r_1 = R_1 \) or the dual parameter pair \( (r_i, \rho_{bc}) \) in the inner boundary condition \( r_1 = r_i \).

The pure neutron Fermi-gas model yields for compact neutron stars a maximum mass of \( M_{\text{max}} = 0.93\, M_{\odot} \) which is in disagreement with observations. Therefore, at least for compact neutron stars, a model of interacting neutron fluid must be used. In 6.2 above we have described a Saxon-Wood-potential model for the nucleon-nucleon interaction, which seems to fit the experiment and the theory in the best way. There will be a critical density (dependent on temperature of course), where a transition from interacting fluid to Fermi-gas takes place, it is plausible to set this density equal to the Saxon-Wood critical density

\[ \rho_c = 0.0417. \]

We made calculations with the TOV-equation using these two models for neutron-based stars and we came to the conclusion that compact neutron stars with mass \( M_0 \leq 3.04\, M_{\odot} \) consist of interacting neutron fluid and neutron shell-stars for \( M_0 \geq 5\, M_{\odot} \) obey the Fermi-gas model. The underlying calculation is the Mathematica-notebook [6].

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This approach yields results, which are described below.

**Neutron stars** consist of interacting neutron fluid and are compact stars with 
\((M_0, R_1) = (0.14, 1.49), \ldots, (3.04, 3.95)\) and the maximum density \(0.048 \leq \rho_{bc} \leq 0.0544 = \rho_{bc\text{max}}\) in sun-units, or shell-stars with \(\rho_{bc} \geq \rho_{bc\text{max}}\) and 
\((M_0, R_1) = (3.04, 3.95), \ldots, (4.91, 4.92)\), neutron star R-M-relation follows approximately a cubic-root-law: \(R \sim M^{1/3}\).

Stellar shell-stars consist of (almost) non-interacting Fermi-gas of neutrons and are thin shell-stars with \(R_i > r_s, R_i = r_s, r_i = r_s\) i.e. the shell is close to the Schwarzschild-radius and its outer edge outside the Schwarzschild-horizon with max. density \(0.0025 \leq \rho_{bc} \leq 0.042\), and obey an almost linear R-M-relation \((M_0, R_i) = (5.5, 9.1), \ldots, (81.3, 91.2)\), independent of \(\rho_{bc}\) for \(\rho_{bc} \geq 0.028\), with redshift factor around 50 for \(M = 80M_{\odot}\).

**Galactic (supermassive) shell-stars** are very thin shell-stars, which obey the **equation-of-state of a white-dwarf** (i.e. gravitation counterbalanced by Fermi-pressure of electron gas) and have an almost linear R-M-relation with redshift factor 20…100.

**Neutron stars**

The parametric solution of the TOV-equation has been carried out for the parameters \((\rho_{bc}, r)\) at the boundary \(r = r_s\) in the range: density \(0.02 \leq \rho_{bc} \leq 0.15\) and inner radius \(0.01 \leq r_i \leq 15\), yielding physical solutions for density \(0.048 \leq \rho_{bc} \leq 0.0544 = \rho_{bc\text{max}}\) and inner radius \(0.01 \leq r_i \leq 3\). The TOV-equation is solved for \(M(r)\) and \(\rho(r)\), and a physical solution is a mathematical solution with \(M \geq 0\) and \(\rho \geq 0\), \(\rho' \leq 0\) and subluminal equation-of-state within a certain interval \(r = \{r_i, r_0\}\), which reaches a point, where \(M'(r) = 0\) and \(\rho(r) = 0\). The radius \(R_i\) and the total mass \(M_0\) is reached at \(M'(R_i) = 0\), the physical solution ends there.

The validity interval for \(\rho\) is explained by the fact, that the sound velocity

\[v_s(\rho) = \frac{\partial P(\rho)}{\partial \rho}\]

must be positive and below 1 (subluminal in c-units).

The parametric mapping of the solutions results in the following dependence for \(M_0(r_s, \rho_{bc}), R_i(r_s, \rho_{bc})\) \((r_s, \rho_{bc}, M_0, R_i\) insun-units) (Figure 8 [6]):

For \(r_i = 0\) the mapping describes the compact neutron stars, resulting in \(R_i(M_0)\) function (Figure 9(a), Figure 9(b), Figure 10 [6]):

The R-M-relation follows approximately a cubic-root-law: \(R \sim M^{1/3}\), with a range of \((M_0, R_i) = (0.14, 1.49), \ldots, (3.04, 3.95)\), i.e. the resulting maximum compact mass is \(M_{\text{max}} = 3.04 M_{\odot}\).

For \(M_0 \geq M_{\text{max}}\) the function \(R_i(t_s = \text{const}, \rho_{bc})\) is flat or slightly decreasing with \(\rho_{bc}\), so one expects the stable configuration to be the one with maximum \(\rho_{bc} = \rho_{bc\text{max}}\) with a range of \((M_0, R_i) = (3.04, 3.95), \ldots, (4.91, 4.92)\).

The admissible mass range ends, where the thickness of the shell above the Schwarzschild-radius becomes very small (minimum 0.01).

So in total the R-M-relation for neutron stars becomes

The maximum mass for a repulsive-hardcore-model for the equation-of-state DD2 [16] is \(2.42M_{\odot}\) from our mapping we have the maximum compact neutron star mass of \(M_{\text{max}} = 3.04M_{\odot}\).
Figure 8. Neutron star: mass $M$ and radius $R$ in dependence of inner boundary density $\rho_{bc}$ and inner radius $r_i$. 

(a)
Figure 9. (a) Compact neutron star: mass-radius $R-M$ relation; (b) Shell neutron star: mass-radius $R-M$ relation.

Figure 10. Neutron star: mass-radius $R-M$ relation.

The actual theoretical limit for neutron star core density is $\rho_{\text{max}} = 3.5 \times 10^{15}$ g/cm$^3 = 0.199$ in sun-units [8] [9].

The limit for $\rho_{\text{bc}}$ reached in our mapping is only $\frac{1}{4}$ of this $\rho_{\text{bc}} = \rho_{\text{bcmax}} = 0.0544$, due to the subluminal-sound-condition and the use of an (attractive) nucleon-nucleon-potential for the nucleon-fluid instead of a pure repulsive-hardcore-model.

The classical argument for the collapse of a neutron star to a black-hole for $\rho_{\text{bc}} > \rho_{\text{max}}$ dating back to Oppenheimer [1], is invalidated here by the simple introduction of shell-star models, where $r_i > 0$, and therefore there is no mass at the center, which means physically, there is only a very diluted nucleon gas there.

Stellar shell-stars (stellar black-holes)
We assume that the underlying equation-of-state state for stellar shell-stars is the Fermi-gas of nucleons with the low-density limit of

\[ P(\rho) = K_i \rho^{5/3}. \]

We make a further plausible assumption that the “edge” of the solution mapping are the physically stable solutions, i.e. the R-M-relation for stellar shell-stars. The edge in this case consists for fixed \( \rho_{bc} < 0.0417 = \rho_{oc} \) of solutions with maximum \( r_i \) (because then the average density in the shell is lowest) and for \( \rho_{bc} = 0.0417 = \rho_{oc} \) it consists of the solutions \((M_0, r_0, R_1)\) at the right boundary \((7 \leq r_i \leq 22, M_0 \geq 4.91 = M_{\text{max}})\), where \( M_{\text{max}} \) is the maximum mass for the neutron star interactive neutron fluid. We assume that at \( \rho_{bc} = \rho_{oc} \) the state transition from the interactive neutron fluid to neutron Fermi gas takes place. This is almost certainly an oversimplification, still we believe that the model describes the reality correctly in principle.

The chosen parameter range of the solution mapping is: density \(0.025 \leq \rho_{bc} \leq 0.0417 = \rho_{oc}\) and inner radius \(0.01 \leq r_i \leq 90\), where \(\rho_{oc}\) is equilibrium value of the nucleon-nucleon-potentials with \(P_{nf}(\rho_{oc}) = 0\), the transition point from the nucleon-fluid to the nucleon-gas phase.

The “edge” of the mapping yields the \((M_0, r_i, R_1)-range of \((M_0, r_0, R_1) = (5.35, 7, 8.49), \ldots, (81.3, 89.6, 91.2)\), where the upper limit is in fact mathematically open, but the “thinning-out” of the solutions for small \(\rho_{bc}\) and large \(r_i\) makes it physically plausible (see Figure 11 [6] below).

The resulting R-M-relation is practically linear and has a maximum mass value of \(M_{\text{max}} = 81.3 M_{\odot}\) (Figure 12(a)).

And the corresponding relative shell thickness \(dR_{\text{rel}} = dR/M\) is (Figure 12(b)) and the relative Schwarzschild-distance \(dR_{\text{rel}} = (R - M)/M\) is (Figure 13)

The inverse of \(dR_{\text{rel}}\) gives roughly the light attenuation factor of \(\{1.7, \ldots, 20.\}\). Taken the attenuation factor and the small relative shell thickness of around 0.02, these stellar shell-stars have approximately the properties expected of a genuine black-hole, when measured from a distance \(r \gg R_i\).
Figure 11. Stellar shell star: mass $M$ and radius $R$ in dependence of inner boundary density $\rho_{bc}$ and inner radius $r_i$.

Figure 12. (a) Stellar shell star: mass-radius $R$-$M$ relation; (b) Stellar shell star: relative shell thickness.
Figure 13. Stellar shell star: relative Schwarzschild-distance.

Entropy of a thin shell star

The celebrated Bekenstein-Hawking formula for the entropy of a black hole reads [17]:
\[ S = \frac{k_B A}{4L_p^2}, \]
where \( A \) is the surface area, \( k_B \) the Boltzmann-constant, and \( L_p \) the Planck-length.

The entropy of a (cold) shell-star with radius \( R \) and thickness \( dR \) in the limit \( R = r_s \), \( dR \ll r_s \), with all particles in the lowest possible energy state, can be easily calculated from the Boltzmann-formula
\[ S = k_B \ln W, \]
where \( W \) is the number of possible micro-states.

With the elementary area \( \pi L_{min}^2 \), where \( L_{min} = L_p \), \( W \) becomes \( W = 2^{\frac{1}{2}(\pi L_{min}^2)} \) (each of the \( N = A \left( \pi L_{min}^2 \right) \) area elements can be occupied or empty), so
\[ S = \frac{k_B A}{\pi L_p^2} \ln 2, \]
which is identical to the Bekenstein-Hawking entropy with the factor \( (\ln 2)/4/\pi = 0.882 \).

Galactic (supermassive) shell-stars

The mean density of a black-hole scales with its radius \( R \) like
\[ \rho(R) = \frac{M}{V} = \frac{R}{\left(\frac{4}{3}\pi R^3\right)} = \frac{3}{4\pi R^2}, \]
i.e. for supermassive black-hole with \( M = 10^6 M_{\odot} \) we have \( \rho \approx 10^{-12} \) in sun units (su).

In the following we use the abbreviation \( MM_{\odot} = 10^6 M_{\odot} \).

The density scale of a white-dwarf star is \( 10^6 \text{ g/cm}^3 = 5.7 \times 10^{-11} \) su [1]. Therefore it is plausible to try a parametric mapping with the white-dwarf equation-of-state, where the underlying Fermi-pressure is that of an electron gas instead of a nucleon gas, i.e. equation-of-state \( P_f(\rho_f) = k_f \rho_f^\gamma \) for a pure Fermi gas, \( \gamma = 5/3 \) if the density is below the critical density \( \rho_c \).

The results for \( M(\rho_f, \rho_c) \), \( R(\rho_f, \rho_c) \) are shown below (Figure 14 [6]).
Figure 14. Galactic shell star: mass $M$ and radius $R$ in dependence of inner boundary density $\rho_{bc}$ and inner radius $r_i$.

From this result one can draw several consequences: first, the actual density is around $10^{-12}$, that is well below the critical density for a white-dwarf of $\rho_c = 0.91 \times 10^6 \text{ g/cm}^3 = 5.17 \times 10^{-11} \text{ su}$: $\gamma = 5/3$ in the equation-of-state is justified. Second, the viable solutions lie to the left of a “ridge” reaching up to masses around $30 M_{\odot}$. Third, a stable solution for a fixed mass will have the highest possible maximum density $\rho_{bc}$ and that will lie on the “ridge”. So one can calculate the $R$-$M$-relation following the “ridge”.

The resulting $R$-$M$-relation is as follows (Figure 15). And the inner radius is (Figure 16).

The $R$-$M$-relation is almost linear, as expected, and goes up to $50 M_{\odot}$. $dR_{rel} = \left( R_i - r_i \right)/M_0$ is the relative thickness (Figure 17), and shows, that the shells are very thin indeed, with a minimum of 0.001. The fourth diagram shows the relative Schwarzschild-distance (Figure 18) $dR_{srel} = \left( R_i - M_0 \right)/M_0$, which has a minimum at $\left[ M_0, dR_{srel} = 0.00142857 \right]$, so that its reciprocal value (approximate light attenuation factor) is around 700. So the overall result is, that the supermassive shell-stars become ever thinner shells, while the distance from the Schwarzschild-horizon is increasing.

7.2. The TOV-Equation: A Case Study for Typical Star Types

In the nearly-rotation-free case the solution of the TOV-equation was calculated for 4 models in sun units with $r_s =$ Schwarzschild radius
Figure 15. Galactic shell star: $R$-$M$-relation.

Figure 16. Galactic shell star: inner radius.

Figure 17. Galactic shell star: relative shell thickness.
Figure 18. Galactic shell star: relative Schwarzschild distance.

\[ r_{ss} = r_1 (\text{sun}) = \frac{2GM_{\text{sun}}}{c^2}, \quad \rho_s = \frac{M_{\text{sun}}}{4\pi r^3/3}, \quad P_s = \rho_s c^2, \]

\[ r_{ss} = 3 \text{ km}, \quad \rho_s = 1.76 \times 10^{16} \text{ g/cm}^3, \quad M_{\text{sun}} = 3 \times 10^{30} \text{ kg}, \]

- average compact neutron star with mass \( M_0 = 0.932M_{\text{sun}} \) radius \( R_1 = 2.767r_{ss} \).
- maximum mass neutron shell-star \( M_0 = 4.91, R = 4.926 \).
- white dwarf with \( M_0 = 0.6M_{\text{sun}} \) radius \( R_1 = 3000r_{ss} \).
- stellar black hole with \( M_0 = 15.69M_{\text{sun}} \) radius \( R_1 = 17.89r_{ss} \), inner radius \( r_i = 17r_{ss} \).
- galactic black hole with \( M_0 = 4.367 \times 10^6M_{\text{sun}} \, R_1 = 4.380 \times 10^6r_{ss}, r_i = 4.356 \times 10^6r_{ss} \).

**Compact neutron star**

parameters = \{k1 = 0.40, gam = 5/3, M0 = 0.932, R1 = 2.76, rhobc = 0.0456, ri = 0.01\};

The mean density is here \( \rho_{\text{mean}} = \frac{M_0}{R_1^3} = 0.04447 \).

The critical density of the neutron Fermi gas with neutron mass \( m_n \) is \( \rho_c = \frac{m_n^4 c^3}{3\pi^2 \hbar^3} = 0.35 \) (see [17]), so the low-density approximation with \( \gamma = 5/3 \) can be used.

**Results TOV:**

density rho (Figure 19), mass \( M \) (Figure 20) are

As can be seen in the \( \rho \)-diagram, the derivative \( \rho' (R_i) = \infty \), because there the equilibrium density \( \rho_c \) with \( P' (r_c) = 0 \) in the pressure is reached.

Maximum mass neutron shell-star

parameters = \{k1 = 0.40, gam = 5/3, M0 = 4.91, R1 = 4.926, rhobc = 0.0544, ri = 3\};

The mean density is here \( \rho_{\text{mean}} = 0.0530 \).

Results TOV:

density rho (Figure 21), mass \( M \) (Figure 22) is

Stellar shell-star (stellar black hole)
parameters = \{k1 = 0.40, \text{gam} = 5/3, M0 = 15.69, R1 = 17.89, rhobc = 0.0359, ri = 17\};

The mean density is here \( \rho_{\text{mean}} = 0.0194 \).

The resulting rho (Figure 23) and \( M \) (Figure 24) are:

Figure 19. Typical compact neutron star: density in dependence of radius.

Figure 20. Typical compact neutron star: mass in dependence of radius.

Figure 21. Shell neutron star: density in dependence of radius.
Figure 22. Shell neutron star: mass in dependence of radius.

Figure 23. Typical stellar shell star: density in dependence of radius.

Figure 24. Typical stellar shell star: mass in dependence of radius.

Here the radius $R_1$ is reached, when $M(r_1 = R_1) = 0$, i.e. $\rho(R_1) = 0$.

White-dwarf star parameters = $\{k_1 = 1.43 \times 10^6, \gamma = 5/3, M_0 = 0.6, R_1 = 3000, \rho_{hbc} = 2.02 \times 10^{-11}, r_i = 0\}$;

The underlying state equation is that of a small-momentum electron Fermi-gas with the critical density $[1]$ \[ \rho_{cw} = \frac{m_i}{3\pi^2\hbar^3} = 0.517 \times 10^{-10} \text{ su}. \]
The mean density is here $\rho_{\text{mean}} = 2.22 \times 10^{-11}$, the maximum deviation of $\rho$ is $\Delta_{\text{max}} \rho = 0.21 \times 10^{-11}$, so the density is practically constant, as expected.

The solution of the TOV-equation yields density $\rho$ (Figure 25), mass $M$ (Figure 26).

Galactic shell-star (supermassive black-hole) parameters = \{k_{11}(\gamma = 5/3) = 0.0243 \times 10^6, k_{12}(\gamma = 4/3) = 0.067 \times 10^4, M_0 = 4.367 \times 10^6, R_1 = 4.380 \times 10^6, \rho_{\text{bc}} = 4.934 \times 10^{-12}, r_i = 4.356 \times 10^6\};

TOV equation was solved with an exterior boundary condition $r_{02} = R_1 (M(r_{02}) = M_0, M'(r_{02}) = 0)$, which is equivalent to the interior boundary condition $r_{01} = r_i (M(r_{01}) = 0, \rho(r_{01}) = \rho_{\text{bc}})$, and with the full Fermi-gas equation-of-state instead of the simple power law $P(\rho) = K_i \rho^n$.

The mean density is here $\rho_{\text{mean}} = 3.16 \times 10^{-12}$.

The “naive” mean density is here $\rho_{\text{mean}} = \frac{M_0}{R_1^3} = 3.16 \times 10^{-12}$, i.e. by a factor 10 lower than the mean density of white dwarf. Therefore, despite its huge mass, the galactic black hole can be described by the state equation of a small-momentum (undercritical) Fermi electron gas with the relative density $x_F = \frac{\rho}{\rho_{\text{cal}}} = 0.0612$ much smaller than that for the white dwarf.

Figure 25. Typical white dwarf star: density in dependence of radius.

Figure 26. Typical white dwarf star: mass in dependence of radius.
TOV-solution for \( \rho \) (in \( 10^{-12} \) units, Figure 27), \( M \) (in \( 10^6 \) units, Figure 28) in \( r \) (in \( 10^6 \) units), is:

Here there is an internal “hole” with a radius \( r_i = 4.356 \times 10^6 \), maximum \( \rho = 4.934 \times 10^{-12} \) at \( r_i \). The inner radius \( r_i \) lies a little below the Schwarzschild-radius \( r_s = M \). The relative shell thickness

**Figure 27.** Typical galactic shell star: density in dependence of radius.

**Figure 28.** Typical galactic shell star: mass in dependence of radius.
\[ dR_{rel} = \left( R_i - r_i \right) / M_0 = 0.00551, \] the relative Schwarzschild-distance
\[ dR_{rel} = \left( R_i - M_0 \right) / M_0 = 0.00290, \] the light attenuation factor is roughly \( 1/dR_{rel} = 344. \)

Furthermore, \( r_i \) is little sensitive to the temperature up to \( T = 10^7 \) K.

As for a stellar black hole, when \( R \) converges to \( r_i = M_0 \), so does the inner radius \( r_i \) and there is no physical solution (with positive \( \rho \) and \( M \)) for a boundary within the horizon.

### 8. The Three Star Models for Kerr-Space-Time with Mass and Rotation

The calculation of Kerr-space-time with mass and rotation was carried out for 3 star models:

- a typical compact neutron star with mass around 1 solar mass;
- a presumably typical stellar shell-star with a mass around 15 solar masses;
- a comparatively small galactic shell-star modelled on the central black-hole in the Milky Way with a mass of around 4 million solar masses.

The angular velocity \( \omega \) was chosen at \( 0.36 \omega_{max} \), i.e. about 1/3 of the maximum value.

We are using the so called “sun units” sun Schwarzschild-radius
\[ r_{ss} = r_s \text{(sun)} = 3 \text{ km}, \] sun mass \( M_s = M \text{(sun)} = 3 \times 10^{30} \text{ kg}, \) sun Schwarzschild-density
\[ \rho_s = \frac{M_s}{4 \pi r_{ss}^3} = 1.76 \times 10^{14} \text{ g/cm}^3, \] sun Schwarzschild-pressure
\[ P_s = \rho_s c^2 = 1.58 \times 10^{35} \text{ J/m}^3 \] for radius \( r \), mass \( M \), density \( \rho \), and pressure \( P \), respectively.

The mass element here is \( M_\theta (r, \theta) d\theta dr \theta \) and the ring mass \( M_\theta (\theta) \) is the differential mass of the \( \theta \)-beam
\[ \frac{d\theta M_\theta (\theta)}{d\theta} = \frac{d\theta M_\theta (r_\theta, \theta)}{d\theta} \text{, the density } \rho \text{ is } 4 \pi \cos (\theta) \rho (r, \theta) d\theta = \partial_r \partial_\theta M_\theta (r, \theta). \]

As for the result values, \( d\theta r_{rel} \) is the maximum relative angular deviation (in \( \theta \)), and error (relative to the test function error): wavefront error is the median (on lattice) algorithm error, the interpolation and the Fourier fit error is the error of the respective fit of the discrete solution on the lattice.

We are using here the \( \theta \)-slicing \( r \)-backward algorithm for the shell-stars. For each of the star models a verification step is run first with the angular velocity \( \omega = 0 \), the result must be the same as in the corresponding TOV-equation. Then a parameter case-study is made for different outer boundary ellipticities \( \Delta R_1 \) at the outer boundary condition in order to find \( \Delta R_1 \) with a minimal mean energy density: this is the stable solution of the Kerr-Einstein equations.

The algorithm yields as the result a pointwise array of values and first and second derivatives of the variables. The variables for the 6 Einstein equations are the 5 Kerr correction-factor functions \( A_0 (r, \theta), \ldots, A_4 (r, \theta) \), and \( \rho (r, \theta) \).

We are using the \( \theta \)-slicing \( r \)-forward algorithm for the compact neutron star. The variables for the 6 Einstein equations are the 5 Kerr correction-factor functions \( A_0 (r, \theta), \ldots, A_4 (r, \theta) \), and \( M_r (r, \theta) \).
The value denotation is:

\[ a = \text{alpha1} \text{ with } a = \frac{J}{Mc} \text{ the angular momentum radius (amr)} \text{ of the Kerr model}, \]
\[ \omega = \text{omega1} \text{ is the angular velocity}, R_i = R_1 = r02 \text{ is the outer radius}, M_0 = M0 \]
\[ r_i \text{ radius variable, th angle variable}, \]
\[ M(r, \theta) = M1(r1, \text{th}) \text{ is the mass function}, A0(r, \theta), \ldots, A4(r, \theta) \text{ Kerr correction-factor functions}, \]
\[ \rho(r, \theta) = \text{rho}(r1, \text{th}) \text{ is the density function}, \]
\[ k_1 \text{ is the parameter in the approximate Fermi-gas equation-of-state} \]
\[ P_1(r_1) = k_1 \rho_1(r_1)^\gamma, \]
\[ \gamma = \text{gam, gam1, gam2 are the exponent}, \]
\[ \text{infac is the moment of inertia factor } \alpha_f, \]
\[ \text{epsi is the singularity cancellation parameter with lim} \text{it}(\text{epsi}) = 0 \text{ introduced to improve the numerical stability in singularities} \]
\[ r_i = \text{riact is the polar inner radius } R_i, \]
\[ \Delta r_i \text{ the inner ellipticity is the difference between the polar } r_i \text{ and the equatorial inner radius } r_{i0}, \]
\[ \Delta R_i \text{ Is the outer ellipticity, with outer radii } R_{i1} = R_i - \Delta R_i \text{ and } R_{i1} = R_i \]
\[ \text{rilow is the minimal radius } r_i \text{ reached in the solution} \]
\[ \rho_{bc} = \text{rhobcx is the boundary condition density} \]
\[ dthrel \text{ is the maximum relative difference of a value dependent on } \theta, \text{ e.g.} \]
\[ dthrel(R_i) = (\text{max}(R_i(\theta)) - \text{min}(R_i(\theta)))/\text{mean}(R_i(\theta)) \]
\[ \text{shell thickness } dR_i = R_i(\theta) - \Delta r_i(\theta) \]

Typical rotating compact neutron star

The underlying star model here is a compact \((r_i = 0)\) neutron star of neutron liquid \((i.e. \text{strongly interacting neutrons}), \text{mass } M_0 = 0.932 \text{ sun-masses, radius } R_i = 2.76 \text{ sun-Schwarzschild-radii } r_{ns} (r_{ns} = 3.0 \text{ km}), \omega = 0.108688. \text{ The underlying calculation is the Mathematica-notebook} [18], \text{the results in} [19].

The full parameters are:
\{alpha1 = 0.331224, omega1 = 0.108688, k1 = 0.4, gam = 5/3., gam1 = 5/3., gam2 = 4/3., M0 = 0.932, dr02rel = 0.33, infac = 2/5., epsi = 0.1, rilow = 0.001, rhobcx = 0.0456\}

The \textit{r-forward} solution is first calculated with the lattice \{nx = 33, ny = 17\} for the rotation-free TOV-case with a TOV-solution as the initial function.

The solution for the Kerr-case starts with this corrected TOV-solution and yields the values:

- outer radius \( R_{i}(\theta) = \{2.83912, \ldots, 2.83722\} \), mean = 2.83897, dthrel = 0.00118.
- total mass \( M02eff = 0.932 \).
- error: \text{med}(\text{err}) = 0.0639 \text{ wavefront}, \text{med}(\text{err}) = 0.0491 \text{ interpolation}, \text{med}(\text{err}) = 0.0492 \text{ Fourier fit.}
- mean energy density = 0.0791.

The resulting density is depicted in \textit{Figure 29}, mass \( M \) in \textit{Figure 30}, and the
effective radius over azimuthal angle $\theta$ in Figure 31.

The density distribution is similar to the TOV-case but with a decrease in $\theta$-direction.

**Figure 29.** Typical rotating compact neutron star: density over $x =$ radius $r_1$, $y =$ angle $\theta$.  

**Figure 30.** Typical rotating compact neutron star: (ring) mass profile for $\theta = 0.1$ (equatorial) and $\theta = 1.4708$ (polar), the outer edge is reached when $M_\theta(r_1) = M_0$.  


The rotation results show very small flattening in the polar direction of \( \delta \theta_{\text{rel}} = 0.00118 \). The neutron star behaves like a fluid because of its “viscosity”, that is, its nuclear interaction and becomes “pumpkin-like”.

Typical rotating stellar shell-star

The star model here is a shell-star \( (r_i > \text{Schwarzschild-radius}) \) with mass \( M_0 = 15.69 \) sun-masses, radius \( R_1 = 17.89 \) sun-Schwarzschild-radii \( r_s (r_s = 3.0 \text{ km}) \), and angular frequency \( \omega = 0.0126 \).

The outer Kerr-horizon is \( r_+ = 15.21 \). The underlying calculation is the Mathematica-notebook [20], the results in [19].

The outer ellipticity \( \Delta R_1 \) is at first a free parameter and calculated from a case-study of minimal mean energy density to \( \Delta R_1 = 0.3 = 0.0168R_1 \).

The full parameters are

\[
\{\alpha_1 = 2.68844, \omega_1 = 0.0126, k_1 = 0.4, R_1 = 17.89, \gamma_1 = 5/-3., \gamma_1 = 5/-3., M_0 = 15.69, \text{infac} = 2/-3., \epsilon_1 = 0.1, r_1 = 15.9, r_{\text{box}} = 0.036\}
\]

The \( r \)-backward solution is first calculated with the lattice \( \{n_x = 17, n_y = 9\} \) for the rotation-free TOV-case as the initial function.

Then a case study with the parameter ellipticity \( \Delta R_1 \) is carried out in order to find the minimal mean energy density, on the set of values \( \Delta R_1 = \{0, -0.3, 0.3, 1., 1.6\} \).

The case study yields a minimum at \( \Delta R_1 = 0.3 \) (cigar-like outer boundary), with mean energy density \( = 0.017004 \).

This energy-minimal solution with \( \Delta R_1 = 0.3 \) yields the values:

- outer radius \( R_1(\theta) = \{17.59, \ldots, 17.89\} \), mean = 17.739, \( \delta \theta_{\text{rel}} = 0.0164 \).
- inner radius \( r_i(\theta) = \{16.704, \ldots, 16.982\} \), mean = 16.823, \( \delta \theta_{\text{rel}} = 0.0162 \).
- total mass \( M_0^{\text{eff}} = 15.69 \), inner boundary \( \max(\rho_{bc}) = 0.035955 \).
- shell thickness \( dR_1 \); mean = 0.916, \( \delta \theta_{\text{rel}} = 0.0448 \).
- mean energy density = 0.0170.

error: med(err) = 0.00238 wavefront, = 0.00573 interpolation, = 0.00786 Fourier fit.

It is interesting to make a comparison with the spherical-outer-boundary so-
with $\Delta R_1 = 0$:
inner radius $r_2(\theta) = \{17.0033, \ldots, 17.0243\}$, mean = 17.020, $d\theta_{rel} = 0.00123$.
total mass $M_{02eff}$ = 15.69, inner boundary $\max(\rho_{bc}) = 0.0375$.
shell thickness $dR_1$: mean = 0.870, $d\theta_{rel} = 0.0241$.
mean energy density = 0.017698.
error: $\text{med}(\text{err}) = 0.00224$ wavefront, = 0.00545 interpolation, = 0.0098 Fourier fit.

The solution with the next higher ellipticity $\Delta R_1 = 1.0$ has the values:
outer radius $r_1(\theta) = \{16.89, \ldots, 17.88\}$, mean = 17.38, $d\theta_{rel} = 0.055$.
inner radius $r_2(\theta) = \{16.01, \ldots, 16.95\}$, mean = 16.438, $d\theta_{rel} = 0.0556$.
total mass $M_{02eff}$ = 15.69, inner boundary $\max(\rho_{bc}) = 0.0359$.
shell thickness $dR_1$: mean = 0.940, $d\theta_{rel} = 0.0894$.
mean energy density = 0.017249.
error: $\text{med}(\text{err}) = 0.00309$ wavefront, = 0.00516 interpolation, = 0.00792 Fourier fit.

The two significant non-spherical features are the relative shell thickness variation $d\theta_{rel}(dR_1)$ and the relative inner ellipticity $d\theta_{rel}(r_1)$. The first depends roughly linearly on the outer ellipticity $\Delta R_1$, plus the value at $\Delta R_1 = 0$ ($d\theta_{rel}(dR_1) = 0.0241$), which is results from rotation. The second, $d\theta_{rel}(r_1)$, is almost equal to the relative outer ellipticity $d\theta_{rel}(R_1)$, plus the small amount at $\Delta R_1 = 0$ ($d\theta_{rel}(R_1) = 0.00123$).

The density distribution is shown in Figures 32-34.

The density distribution increases in $\theta$-direction ($\theta = 0$ is equatorial, $\theta = \pi/2$ axial).

The mass distribution is shown in Figure 35, Figure 36.

The physical mass distribution ends at the inner boundary at $r_1 = 16.7$, where the density jumps to $\rho = 0$.

Figure 32. Typical rotating stellar shell star: density over $x =$ radius $r_1$, $y =$ angle $\theta$. 
A remarkable result, distinct from the case of the neutron star, is the shape with rotation. The energy-minimal stellar shell-star behaves like a ball of neutron gas (negligible interaction) and decreases slightly its equatorial radius, so that, speaking naively, the increased gravitation counteracts the centrifugal force, the shell-star becomes “cigar-like”, with the shell thickness approximately constant.
Figure 36. Typical rotating stellar shell star: Fourier-fitted total mass in dependence of radius.

The stellar shell-star has all its mass concentrated within a thin shell \((dR_1 = 0.916)\) which is situated outside its Schwarzschild-radius \(M_0 = 15.69\), where the minimum distance from the horizon is 
\[\min(R_1(\theta)) - M_0 = 1.01,\]
therefore the light energy loss is approximately (in Newtonian approximation) 
\[M_0/\min(R_1(\theta)) = 0.9395\]
and the attenuation factor \(1/(1 - M_0/\min(R_1(\theta))) = 16.53\), it means that visible green light of 0.514 \(\mu\)m is shifted to 8.50 \(\mu\)m into middle-range infrared.

Typical rotating galactic shell-star

This is modelled (approximately) on the central black-hole in the Milky Way with mass \(M_0 = 4.36\) mega-sun-masses \((MM)\), radius \(R_1 = 4.38\) mega-sun-Schwarzschild-radii \((13.14 \times 10^6\) km, \(Mrss)\) [21].

The underlying calculation is the Mathematica-notebook [22], the results in [23].

The outer Kerr-horizon is \(r_+ = 4.26 Mrss\).

In order to maintain numerical performance, we are using for mass and distance \(10^6\) (mega) units \(10^6 Ms\) and \(10^6 rss\) and for density \(10^{-12}\) (mega\(^{-2}\)) unit \(10^{-12} \rho_s\).

Like in the case of the stellar shell-star, the outer ellipticity \(\Delta R_1\) is at first a free parameter and calculated from a case-study of minimal mean energy density to \(\Delta R_1 = -2 \text{ dTOV}\), where dTOV is the shell thickness of the spherical shell-star dTOV = 0.057.

The full parameters are:
\[
\{\alpha_1 = 0.670047, \omega_1 = 0.05239, k_1 = 0.0243, k_2 = 0.067, R_1 = 4.38, \gamma = 5/3, \\
\gamma_1 = 5/3, \gamma_2 = 4/3, M_0 = 4.36731, \text{ infac = 2/3, epsi = 0.0024, rilow = 4.3, rhobcx = 4.915}\}
\]

The \(r\)-backward solution is first calculated with the lattice \(\{n_x = 17, n_y = 9\}\) for the rotation-free TOV-case as the initial function.

Then a case study with the parameter ellipticity \(\Delta R_1\) is carried out in order to find the minimal mean energy density, on the set of values \(\Delta R_1 = \{0.1, -0.3, -1, -2\} \text{ dTOV}\).

The case study yields a minimum at \(\Delta R_1 = -2 \text{ dTOV} = -0.114\) (pancake-like
outer boundary), with a mean energy density $= 1.734$.

This energy-minimal solution yields the values:

- outer radius $R(\theta) = \{4.494, \ldots, 4.38\}$, mean $= 4.43654$, dthrel $= 0.025367$.
- inner radius $r(\theta) = \{4.46456, \ldots, 4.34109\}$, mean $= 4.40029$, dthrel $= 0.02765$.
- total mass $M_{02\text{eff}} = 4.3673$, inner boundary max$(\rho_{bc}) = 4.96075$.
- shell thickness $dR$: mean $= 0.035256$, dthrel $= 0.2507$.
- mean energy density $= 1.73479$.
- error: med$(\text{err}) = 0.0122$ wavefront, $= 0.000404$ interpolation, $= 0.000437$ Fourier fit.

In comparison, the spherical-outer-boundary solution with $\Delta R = 0$ yields the values:

- inner radius $r(\theta) = \{4.34861, \ldots, 4.3423\}$, mean $= 4.343$, dthrel $= 0.00145$.
- total mass $M_{02\text{eff}} = 4.36731$, inner boundary max$(\rho_{bc}) = 4.68324$.
- shell thickness $dR$: mean $= 0.0370$, dthrel $= 0.1736$.
- mean energy density $= 1.75914$.
- error: med$(\text{err}) = 0.0101$ wavefront, $= 0.000390$ interpolation, $= 0.000324$ Fourier fit.

In contrast to the stellar shell-star, here the relative variation of the shell thickness for the spherical-outer-boundary solution is smaller by a factor of 20 as compared to the minimal solution with a high outer ellipticity, so here there is a dependence of the shell thickness on the ellipticity.

The density distribution is shown in Figures 37-39.

The density distribution increases in $\theta$-direction.

The mass distribution is shown in Figure 40, Figure 41.

The physical mass distribution ends at the inner boundary at $r_i = 4.46456$, where the density jumps to $\rho = 0$. The fit extrapolates it to lower $r$-values.

The maximum distance from the horizon is $\max(r_{02\text{eff}}) - r = 0.125$, therefore the minimal light energy attenuation is roughly $4.262/0.125 = 34$, it means that visible green light of 0.514 μm is shifted to 17 μm into far-infrared.

![Figure 37. Typical rotating galactic shell star: density over $x = \text{radius } r$, $y = \text{angle } \theta$, radius in mega-$r_{(\text{sun})}=3 \times 10^6$ km, $\theta$ in radian.](image)
The galactic shell-star has all its mass concentrated within a thin shell 
\((dR_i = 0.0362)\) which has its inner radius inside and its outer radius outside its 
Schwarzschild-radius \(M_0 = 4.367\), where the minimum distance from the horizon is 
\(\min(R_i(\theta)) - M_0 = 0.0127\), 
therefore the light energy loss is approximately (in Newtonian approximation)
Figure 41. Typical rotating galactic shell star: total mass $M_1(r_1)$ over radius, radius in $r_{\text{sun}}=3$ km, $\theta$ in radian.

Figure 42. Mass distribution of stellar black-holes and neutron stars measured by LIGO.

$M_0/\min(R_1(\theta)) = 0.9971$ and the attenuation factor $1/(1 - M_0/\min(R_1(\theta))) = 345$, it means that x-ray-radiation from in-falling matter from the accretion disc with an energy of 5 keV and $\lambda = 0.2$ nm is shifted to $\lambda = 69$ nm, that is into hard UV-radiation.

9. Experimental Evidence with Recent LIGO and X-Ray Measurements

In November 2018, the LIGO cooperation published the newest statistics of neutron stars and black holes, based on gravitational waves and x-ray measurements [24].

The resulting mass distribution for black-holes and neutron stars is shown in Figure 42 [24].

From these results, we can deduce a confirmed mass range for neutron stars of $0.9 M_{\text{sun}} \leq M_n \leq 2.9 M_{\text{sun}}$ and for stellar black holes $5M_{\text{sun}} \leq M_{\text{bh}} \leq 80 M_{\text{sun}}$.

This mass range confirms our results from chapter 7 for compact neutron stars in the range $M \leq 3.04 M_{\text{sun}}$, but not for the neutron shell stars in the range $3.04 M_{\text{sun}} \leq M \leq 4.91 M_{\text{sun}}$. 

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It means that the fluid (non-compact) shell stars are not stable, and only Fermi-gas shell stars are stable.

The mass range for stellar shell star black holes deduced in chapter 7: 
\[5.5M_{\odot} \leq M_{bh} \leq 81.3M_{\odot}\] is confirmed, considering that the uncertainty of mass measurement via x-rays is easily \(\delta M_{bh} \geq 0.5M_{\odot}\).

10. Conclusions

We introduce in chap. 6 an eos for the nucleon-fluid in the density range \(\rho_c \leq \rho \leq \rho_m\) where \(\rho_c = 0.0417\rho_s\) and \(\rho_m = 0.0544\rho_s\) (sun units \(r_s = 3\) km, \(\rho_s = 1.76 \times 10^{16}\) g/cm\(^3\)), which is based on measurement data for the nucleon-nucleon-potential. This suggests, that there is a phase transition at \(\rho = \rho_c\) from the (interacting) nucleon-fluid to the (weakly interacting) nucleon Fermi-gas.

Based on these 2 eos’s the results for the TOV-equation in chap. 7 are as follows.

Neutron stars obey the nucleon fluid eos and there are compact neutron stars in the range \((M_0, R_1) = (0.14M_{\odot}, 1.49r_s), \ldots, (3.04M_{\odot}, 3.95r_s)\), the \(R-M\)-relation follows approximately a cubic-root-law: \(R \sim M^{1/3}\).

Neutron shell-stars in the range \((M_0, R_1) = (3.04M_{\odot}, 3.95r_s), \ldots, (4.91M_{\odot}, 4.92r_s)\) are not stable.

Stellar shell-stars exist in the range of \((M_0, R_1) = (5.5M_{\odot}, 9.1r_s), \ldots, (81.3M_{\odot}, 91.2r_s)\).

The underlying equation-of-state is the Fermi-gas of nucleons with the eos \(P(\rho) = K_{e}\rho^{5/3}\). The resulting \(R-M\)-relation is practically linear and has a maximum mass value of \(M_{max} = 81.3M_{\odot}\). The light attenuation factor (redshift) is roughly \(\{1.7, \ldots, 50\}\). Taken the redshift and the small relative shell thickness of around 0.02, these stellar shell-stars have approximately the properties expected of a genuine black-hole, when measured from a distance \(r \gg R_1\). Furthermore, the phase space volume of a thin spherical shell is proportional to its surface \(A\), which approximates the Bekenstein-Hawking black-hole entropy formula \(S = (k_B / L_p) A / 4\).

The galactic (supermassive) shell-stars have the density scale and the eos of a white-dwarf-star, i.e. of an electron Fermi-gas. The \(R-M\)-relation is almost linear and goes from \(1M_{\odot}\) up to \(50M_{\odot}\) \((M_{\odot} = 10^8M_{\odot}, M_{\odot} = 10^6r_s)\). \(dR_{rel} = (R_1 - r) / M_0\) is the relative thickness, and shows, that the shells are very thin indeed, with a minimum of 0.001. The relative Schwarzschild-distance \(dR_{rel} = (R_1 - M_0) / M_0\) has a minimum at \(\{M_0, dR_{rel}\} = \{7M_{\odot}, 0.00142857\}\), the redshift is around 700. So the overall result is, that the supermassive shell-stars become even thinner shells, while the distance from the Schwarzschild-horizon is increasing.

In chap. 8 we present numerical results for rotating stars of the 3 types compact neutron star, stellar shell-star and galactic shell-star.

The angular velocity \(\omega\) was chosen at \(\omega = 0.36\omega_{max}\), i.e. about 1/3 of the maximum.
The compact neutron star with $M_0 = 0.932M_{\odot}$, $R_{1y} = 2.8372r_s$, $R_{1x} = 2.8391r_s$, $\omega = 0.1087$, has the relative ellipticity of $d\theta_{\text{rel}} = 0.00118$. The neutron star behaves like a fluid because of its “viscosity”, that is, its nuclear interaction, and becomes slightly “pumpkin-like”.

The stellar shell-star with $M_0 = 15.74M_{\odot}$, $R_{\text{mean}} = 17.74r_s$, $\omega = 0.0126$, has maximum density $\rho_{bc} = 0.03595\rho_s$, outer radii $R_{1y} = 17.89r_s$, $53.67$ km, $R_{1x} = 17.59r_s$, inner radii $r_{iy} = 16.98r_s$, $r_{ix} = 16.70r_s$, inner rel. ellipticity $d\theta_{\text{rel}} = 0.0162$. The redshift is 16.53.

The stellar shell-star behaves like a ball of neutron gas (negligible interaction) and decreases slightly its equatorial radius, so that, speaking naively, the increased gravitation counteracts the centrifugal force, the shell-star becomes “cigar-like”, with the shell thickness approximately constant.

The galactic shell-star is modelled (approximately) on the central black-hole in the Milky Way with mass $M_0 = 4.368M_{\odot}$, $M_{\text{sun}} = 10^8M_{\odot}$, $M_{\text{rs}} = 10^8r_s$, radius $R = 4.38Mr_s (=13.14 \times 10^6$ km), $\omega = 0.05239$.

It has maximum density $\rho_{bc} = 4.961 \times 10^{-12}\rho_s$, outer radii $R_{1y} = 4.38Mr_s$, $R_{1x} = 4.494Mr_s$, inner radii $r_{iy} = 4.341Mr_s$, $r_{ix} = 4.464Mr_s$, inner rel. ellipticity $d\theta_{\text{rel}} = 0.0276$.

The redshift is roughly 345. The galactic shell-star is a shell object with a thin mass shell ($\Delta R = 0.0352Mr_s$) situated close above its outer Kerr horizon $r_+ = 4.26Mr_s$. The polar radius is smaller than the equatorial radius, so the outer shape and the inner shape are both pancake-like.

The overall result is, that the introduction of numerical shell-star solutions of the TOV- and Kerr-Einstein-equations creates shell-star star models, which mimic closely the behaviour of abstract black holes and satisfy the Bekenstein-Hawking entropy formula, but have finite redshifts and escape velocity $v < c$, no singularity, no information loss paradox, and are classical objects, which need no recourse to quantum gravity to explain their behaviour.

**Conflicts of Interest**

The author declares no conflicts of interest regarding the publication of this paper.

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