Abstract. We study gravitational collapse of uniformly rotating neutron stars to Kerr black holes, using a new three-dimensional, fully general relativistic hydrodynamics code, which uses high-resolution shock-capturing techniques and a conformal traceless formulation of the Einstein equations. We investigate the gravitational collapse by carefully studying not only the dynamics of the matter, but also that of the trapped surfaces, i.e. of both the apparent and event horizons formed during the collapse. The use of these surfaces, together with the dynamical horizon framework, allows for a precise measurement of the black-hole mass and spin. The ability to successfully perform these simulations for sufficiently long times relies on excising a region of the computational domain which includes the singularity and is within the apparent horizon. The dynamics of the collapsing matter is strongly influenced by the initial amount of angular momentum in the progenitor star and, for initial models with sufficiently high angular velocities, the collapse can lead to the formation of an unstable disc in differential rotation.

1. Introduction

The numerical investigation of gravitational collapse of rotating stellar configurations leading to black-hole formation is a long standing problem in numerical relativity. However, it is through numerical simulations in general relativity that one can hope to improve our knowledge of fundamental aspects of Einstein’s theory such as the cosmic censorship hypothesis and black-hole no-hair theorems, along with that of current open issues in relativistic astrophysics research, such as the mechanism responsible for gamma-ray bursts. Furthermore, numerical simulations of stellar gravitational collapse to black holes provide a unique mean of computing the gravitational waveforms emitted in such events, believed to be among the most important sources of detectable gravitational radiation.

However, the modelling of black-hole spacetimes with collapsing matter-sources in multidimensions is one of the most formidable efforts of numerical relativity. The ability to perform long-term numerical simulations of self-gravitating systems in general relativity strongly
depends on the formulation adopted for the Einstein equations. Building on the experience developed with lower-dimensional formulations, Nakamura, Oohara and Kojima \cite{1} presented in 1987 a conformal traceless reformulation of the ADM system which subsequent authors (see, e.g., \cite{2, 3, 4, 5, 6, 7, 8, 9}) gradually showed to be robust enough to accomplish such a goal for different classes of spacetimes, including black holes and neutron stars (both isolated and in coalescing binary systems). The most widespread version developed from this formalism, which we refer to here as the NOK formulation, was given by \cite{2, 3} and is commonly referred to as the BSSN formulation.

In addition to the improvements achieved in the formulation of the field equations, successful long-term 3D evolutions of black holes in vacuum have been obtained in the last few years using excision techniques (see, e.g., \cite{10, 11, 12, 13, 14, 15, 16, 17, 18, 19}), although the original idea is much older \cite{20}. In this approach, the spacetime region within the black-hole horizon is causally disconnected and so can be safely ignored without affecting the evolution outside the horizon as long as suitable boundary conditions are specified at the excision surface. The simulations presented here show the applicability of excision techniques also in non-vacuum spacetimes, namely during the collapse of rotating neutron stars to Kerr black holes (for details of the method, see \cite{21}).

The presence of rotation in the collapsing stellar models requires multidimensional investigations, either in axisymmetry or in full 3D. The numerical investigations of black-hole formation (beyond spherical symmetry) started in the early 1980’s with the pioneering work of Nakamura \cite{22}. He adopted the (2+1)+1 formulation of the Einstein equations in cylindrical coordinates and introduced regularity conditions to avoid divergences at coordinate singularities. In a series of papers \cite{23, 24, 25, 26}, Bardeen, Stark and Piran studied the collapse of rotating relativistic polytropes to black holes, succeeding in computing the associated gravitational radiation. The initial model was a spherically symmetric relativistic polytrope of mass $M$ in equilibrium. The gravitational collapse was induced by lowering the pressure in the initial model by a prescribed (and often very large) fraction and by simultaneously adding an angular momentum distribution approximating rigid-body rotation. More recently, Shibata \cite{5, 27} has performed axisymmetric simulations of the collapse of rotating supramassive neutron stars to black holes finding an upper limit to the mass of a possible disc as being less than $10^{-3}$ of the initial stellar mass. Three-dimensional, fully-relativistic simulations of the collapse of supramassive uniformly rotating neutron stars to rotating black holes were presented in \cite{4} showing no evidence of massive disc formation or outflows, in agreement new simulations reported by \cite{28}.

Here, we present the results of new, fully 3D simulations of gravitational collapse of uniformly rotating neutron stars. For the first time in such 3D simulations, we have detected the event horizon of the forming black hole and showed that it can be used to achieve a more accurate determination of the black-hole mass and spin than it would be otherwise possible using the area of the apparent horizon. We have also considered several other approaches to measure the properties of the newly formed Kerr black hole, including the recently proposed isolated and dynamical horizon frameworks. A comparison among the different methods indicates that the dynamical horizon approach is simple to implement and yields estimates which are accurate and more robust than those of the equivalent methods. Our results also show that the dynamics of the collapsing matter is strongly influenced by the initial amount of angular momentum in the progenitor neutron star, which, when sufficiently high, leads to the formation of an unstable flattened disc.

The simulations are performed with a new general relativistic hydrodynamics code, using a Cartesian grid (a first description of the code was given in \cite{29}). The code incorporates the expertise developed over the past few years in the numerical solution of the Einstein equations and of the hydrodynamics equations in a curved spacetime (see \cite{6, 7}, but also \cite{30} and references
2. Unstable Equilibrium Models

As mentioned earlier, this study is specially dedicated to the gravitational collapse of slowly and rapidly rotating supramassive relativistic stars, in uniform rotation, that have become unstable to axisymmetric perturbations. Given equilibrium models of gravitational mass $M$ and central energy density $e_c$ along a sequence of fixed angular momentum or fixed rest mass, the Friedman, Ipser & Sorkin criterion $\partial M/\partial e_c = 0$ [32] can be used to locate the exact onset of the secular instability to axisymmetric collapse. The onset of the dynamical instability to collapse is located near that of the secular instability but at somewhat larger central energy densities. Unfortunately, no simple criterion exists to determine this location, but the expectation mentioned above has been confirmed by the simulations performed here and by those discussed in [4]. Note that in the absence of viscosity or strong magnetic fields, the star is not constrained to rotate uniformly after the onset of the secular instability and could develop differential rotation. In a realistic neutron star, however, viscosity or intense magnetic fields are likely to enforce a uniform rotation and cause the star to collapse soon after it passes the secular instability limit.

The initial data for our simulations are constructed using a 2D numerical code, that computes accurate stationary equilibrium solutions for axisymmetric and rapidly rotating relativistic stars in polar coordinates [33]. The data are then transformed to Cartesian coordinates using standard coordinate transformations. The same initial data routines have been used in previous 3D simulations [6, 7, 34] and details on the accuracy of the code can be found in [35].

For simplicity, we have focused on initial models constructed with the usual relativistic
Table 1. Equilibrium properties of the initial stellar models. The different columns refer respectively to: the central rest-mass density \( \rho_c \), the ratio of the polar to equatorial coordinate radii \( r_p/r_e \), the gravitational mass \( M \), the circumferential equatorial radius \( R_e \), the angular velocity \( \Omega \), the ratio \( J/M^2 \) where \( J \) is the angular momentum, the ratio of rotational kinetic energy to gravitational binding energy \( T/|W| \), and the “height” of the corotating ISCO \( h_+ \). All models have been computed with a polytropic EOS with \( K = 100 \) and \( \Gamma = 2 \).

| Model | \( \rho_c \times 10^{-3} \) | \( r_p/r_e \) | \( M \) | \( R_e \) | \( \Omega \times 10^{-2} \) | \( J/M^2 \times 10^{-2} \) | \( T/|W| \) | \( h_+ \) |
|-------|----------------|-------------|--------|-------|----------------|----------------------------|----------------|---------|
| S1    | \( 3.154 \)   | \( 0.95 \)  | \( 1.666 \) | \( 7.82 \) | \( 1.69 \)   | \( 0.207 \) | \( 1.16 \) | \( 1.18 \) |
| S2    | \( 3.066 \)   | \( 0.85 \)  | \( 1.729 \) | \( 8.30 \) | \( 2.83 \)   | \( 0.363 \) | \( 3.53 \) | \( 0.51 \) |
| S3    | \( 3.013 \)   | \( 0.75 \)  | \( 1.798 \) | \( 8.90 \) | \( 3.49 \)   | \( 0.470 \) | \( 5.82 \) | \( 0.04 \) |
| S4    | \( 2.995 \)   | \( 0.65 \)  | \( 1.863 \) | \( 9.76 \) | \( 3.88 \)   | \( 0.545 \) | \( 7.72 \) | \( - \) |
| D1    | \( 3.280 \)   | \( 0.95 \)  | \( 1.665 \) | \( 7.74 \) | \( 1.73 \)   | \( 0.206 \) | \( 1.16 \) | \( 1.26 \) |
| D2    | \( 3.189 \)   | \( 0.85 \)  | \( 1.728 \) | \( 8.21 \) | \( 2.88 \)   | \( 0.362 \) | \( 3.52 \) | \( 0.58 \) |
| D3    | \( 3.134 \)   | \( 0.75 \)  | \( 1.797 \) | \( 8.80 \) | \( 3.55 \)   | \( 0.468 \) | \( 5.79 \) | \( 0.10 \) |
| D4    | \( 3.116 \)   | \( 0.65 \)  | \( 1.861 \) | \( 9.65 \) | \( 3.95 \)   | \( 0.543 \) | \( 7.67 \) | \( - \) |

polytropic EOS, choosing \( \Gamma = 2 \) and polytropic constant \( K = 100 \) to produce stellar models that are, at least qualitatively, representative of what is expected from observations of neutron stars.\(^9\) More specifically, we have selected four models located on the line defining the onset of the secular instability and having polar-to-equatorial axes ratio of roughly 0.95, 0.85, 0.75 and 0.65 (these models are indicated as S1–S4 in Fig.1), respectively. Four additional models were defined by increasing the central energy density of the secularly unstable models by 5%, keeping the same axis ratio. These models (indicated as D1–D4 in Fig.1) were expected and have been found to be dynamically unstable. Figure 1 shows the gravitational mass as a function of the central energy density for equilibrium models constructed with the chosen polytropic EOS. Table 1 summarizes the main equilibrium properties of the initial models.

3. Dynamics of the Matter
The slowly rotating model D1 collapses nearly spherically symmetric. However, for the most rapidly rotating model D4, the large angular initial velocity produces significant deviations from a spherical infall, as the collapse proceeds. Indeed, the parts of the star around the rotation axis that are experiencing smaller centrifugal forces collapse more promptly and, as a result, the configuration increases its oblateness. At about \( t = 0.64 \) ms an apparent horizon appears. Soon after this, the central regions of the computational domain are excised, preventing the code from crashing. An unstable disc of low-density matter remains near the equatorial plane, orbiting at very high velocities \( > 0.2 \) \( c \), see Figure 2. By a time \( t = 0.85 \) ms, essentially all (i.e. more than 99.9\%) of the residual stellar matter has fallen within the trapped surface of the apparent horizon and the black hole thus formed approaches the Kerr solution.

4. Dynamics of the Horizons
In order to investigate the formation of a black hole in our simulations, we have used horizon finders, available through the Cactus framework, which compute both the apparent horizon and

\(^9\) Unless otherwise noted, we use dimensionless quantities by setting \( c = G = M_\odot = 1 \).
Figure 2. Magnified view of the final stages of the collapse of our most rapidly rotating model. Note that a region around the singularity that has formed is excised from the computational domain and this is indicated as an area filled with squares. Also shown with a thick dashed line is the coordinate location of the apparent horizon. An unstable disk forms during collapse.

The event horizon. The apparent horizon, is calculated at every time step. In contrast, the event horizon, is computed a posteriori, once the simulation is finished, by reconstructing the full spacetime from the 3D data each simulation produces. In all cases considered, we have found that the event horizon rapidly grows to its asymptotic value after formation. With a temporal gap of \( \sim 10M \) after the formation of the event horizon, the apparent horizon is found and then it rapidly approaches the event horizon, always remaining within it.

We measure the mass of the newly formed black hole and compare it to the ADM mass of the spacetime computed by the initial data solver. The first and simplest method of approximating the black-hole mass is to note that, for a Kerr (or Schwarzschild) black hole, the mass can be found directly in terms of the event-horizon geometry as

\[
M = \frac{C_{\text{eq}}}{4\pi},
\]

where \( C_{\text{eq}} \equiv \int_0^{2\pi} \sqrt{g_{\phi\phi}} d\phi \) is the proper equatorial circumference, see Figure 3.

Following other approaches, a major difficulty in an accurate measurement of \( M \) lies in the calculation of its non-irreducible part, i.e. in the part that is proportional to the black-hole angular momentum \( J \). We now discuss a number of different ways to calculate \( J \) from the present simulations; these measurements will then be used to obtain alternative estimates of \( M \).

One idea is to look at the distortion of the horizon using the ratio of polar and equatorial proper circumferences, \( C_r \equiv C_{\text{pol}}/C_{\text{eq}} \). For a perturbed Kerr black hole this is expected to oscillate around the asymptotic Kerr value with the form of a quasi-normal mode (QNM). By fitting to this mode we extract an estimate of the angular momentum parameter \( a/M_{\text{hor}} \) from the relation [36]

\[
\frac{a}{M_{\text{hor}}} = \sqrt{1 - (-1.55 + 2.55C_r)^2},
\]
Figure 3. Evolution of the event-horizon mass $M = C_{\text{eq}}/4\pi$ for models D1 and D4. Different lines refer to the different initial guesses for the null surface and are numbered 0, 1 and 2. Note that they converge exponentially to the correct event-horizon surface for decreasing times. The differences at late times introduces a systematic error in evaluating the mass. Shown with the horizontal lines are the ADM masses in two cases and the lower limit on the error-bar as measured from the initial data.

where we have indicated with $M_{\text{hor}}$ the black-hole mass as measured from expression (2), which coincides with $M$ only if the spacetime has become axisymmetric and stationary. The fit through expression (2) is expected to be accurate to $\sim 2.5\%$ [36].

The fit itself depends on an initial guess for $a/M_{\text{hor}}$ and we start from a Schwarzschild black hole and iterate till the desired convergence is reached. This measure is not gauge invariant, although Equation (2) is independent of the spatial coordinates up to the definition of the circumferential planes, but works adequately with the gauges used here. The fit is best performed shortly after black-hole formation as the oscillations are rapidly damped. This minimizes numerical errors but in those cases where matter continues to be accreted, it may lead to inaccurate estimates of the angular momentum. Using expression (2) to estimate the angular momentum $J$ introduces an error, if the black hole has not yet settled to a Kerr solution. Having this in mind, however, it is possible to estimate the angular momentum as

$$J = \left(\frac{a}{M_{\text{hor}}}\right) M_{\text{hor}} M \simeq \left(\frac{a}{M_{\text{hor}}}\right) M^2.$$

A second method of approximating $J$ and hence measuring $M$ is to use the isolated and dynamical horizon frameworks of Ashtekar and collaborators [37, 38, 39, 40, 41]. This assumes the existence of an axisymmetric Killing vector field intrinsic to a marginally trapped surface, such as an apparent horizon. This gives an unambiguous definition of the spin of the black hole and hence of its total mass. If there is an energy flux across the horizon, the isolated horizon framework needs to be extended to the dynamical horizon formalism [40, 42].

In practice, our approach to the dynamical horizon framework has been through the use of a code by Schnetter which implements the algorithm of [41] to calculate the horizon quantities.
The advantage of the dynamical horizon framework is that it gives a measure of mass and angular momentum which is accurately computed locally, without a global reconstruction of the spacetime. In the case of a slowly rotating model the estimate from the dynamical horizon finder is perfectly stable, indicating that an approximately stationary Kerr black hole has been formed by the time the simulation is terminated. In the case of a rapidly rotating model, however, this is no longer the case as matter continues to be accreted also at later times, when the errors have also increased considerably. As a result, the measure of the spin through the dynamical horizon finder is less accurate and does not seem to have stabilized by the time the simulation ends. This may indicate that the final black hole has not settled down to a Kerr black hole on the timescales considered here.

A third method for computing \( J \) only applies if an event horizon is found and if its angular velocity has been measured. In a Kerr background, in fact, the generators of the event horizon rotate with a constant angular velocity \( \omega \equiv -g_{\phi\phi}/g_{\phi\phi} = \sqrt{g_{tt}/g_{\phi\phi}} \). In this case the generator velocity can be directly related to the angular momentum parameter as

\[
\frac{a}{M} = \frac{J}{M^2} = \left[ \frac{A\omega^2}{\pi} \left(1 - \frac{4\omega^2}{4\pi}\right) \right]^{1/2}.
\] (4)

As with the previous approximations, expression (4) is strictly valid only for a Kerr black hole and will therefore contain a systematic error which, however, decays rapidly as the black-hole perturbations are damped. On the other hand, the event horizon generator velocities have the considerable advantage that everything is measured instantaneously and the values of \( \omega \) are valid whether or not the background is an isolated Kerr black hole.

It was shown by Christodoulou that, in the axisymmetric and stationary spacetime of a Kerr black hole, the square of the black-hole mass \( M \) is given by

\[
M^2 = M_{\text{irr}}^2 + \frac{4\pi J^2}{A} = \frac{A}{16\pi} + \frac{4\pi J^2}{A},
\] (5)

where \( M_{\text{irr}} \) is the irreducible mass, \( A \) is the event-horizon proper area, and \( J \) is the black-hole angular momentum. As the black hole approaches a stationary state at late times, the apparent and event horizons will tend to coincide and in that case the mass of the black hole is very well approximated by the above formula.

We have applied the above formula, using the various methods for measuring the angular momentum \( J \). In particular, using the method for obtaining \( J \) from the distortion of the event horizon, through Equation (3), the black-hole mass is given by

\[
M^2 = \frac{A}{8\pi} \left( \frac{M_{\text{hor}}}{a} \right)^2 \left[ 1 - \sqrt{1 - \left( \frac{a}{M_{\text{hor}}} \right)^2} \right].
\] (6)

If, on the other hand, \( J \) is found from the angular velocity \( \omega \) of the event horizon, then it is possible to use (4) in (5) and obtain

\[
M^2 = \frac{A}{16\pi - 4A\omega^2}.
\] (7)

In the framework of dynamical horizons, expression (5) holds for any axisymmetric isolated or dynamical horizon, independently of whether it is stationary.

For the case of the slowly rotating model D1, all of the matter rapidly collapses into the black hole, and the different estimates of the total mass agree very well. However, already in this slowly
rotating case the irreducible mass of the apparent horizon is noticeably lower. For the rapidly rotating model, the contribution from the spin energy is considerably larger and noticeable differences appear among the different approaches. Since all seem to have systematic errors, this makes it less trivial to establish which method to prefer. On one hand, those methods using information from the event-horizon equatorial circumference or that fit the perturbations of the event horizon \[\text{i.e. Equations (1) and (6)}\] seem to provide accurate estimates at earlier times but suffer of the overall inaccuracy at later stages, when the initial guesses for the null surface are distinct. On the other hand, those methods that measure the angular velocity of the null generators \[\text{i.e. Equation (7)}\] or that use the dynamical horizon framework, produce reasonably accurate estimates, that converge with resolution, that monotonically grow in time and that are within the error-bar of the initial estimate of $M_{\text{ADM}}$. Furthermore, in the case of the dynamical horizon framework, this is not only physically expected, given that a small but non-zero fraction of the matter continues to be accreted nearly until the end of the simulation, but it is also guaranteed analytically. Although the direct comparison of many different methods employed here has provided valuable information on the dynamics of the system, we have found the dynamical horizon framework to be simple to implement, accurate and not particularly affected by the errors from which equivalent approaches seem to suffer. As a result, we recommend its use as a standard tool in numerical relativity simulations.

5. Reconstructing the global spacetime

All of the results presented and discussed in the previous Section describe only a small portion of the spacetime which has been solved during the collapse. In addition to this, it is interesting and instructive to collect all of these pieces of information into a \textit{global} description of the spacetime and look for those features which mark the difference between the collapse of slowly and rapidly rotating stellar models. As we discuss below, these features emerge in a very transparent way within a global view of the spacetime.

To construct this view, we use the worldlines of the most representative surfaces during the collapse, namely those of the equatorial stellar surface, of the apparent horizon and of the event horizon. For all of them we need to use properly defined quantities and, in particular, circumferential radii. The results of this spacetime reconstruction are shown in Fig. 4, which refers to the collapse of model D4. The different lines indicate the worldlines of the circumferential radius of the stellar surface (dotted line), as well as of the apparent horizon (dashed line) and of the event horizon (solid line). Note that for the horizons we show both the equatorial and the polar circumferential radii, with the latter being always smaller than the former. For the stellar surface, on the other hand, we show the equatorial circumferential radius only. This is because the calculation of the stellar polar circumferential radius requires a line integral along the stellar surface on a given polar slice. Along this contour, one must use a line element which is suitably fitted to the stellar surface and diagonalized (see [44] for a detailed discussion). In the case of model D4, however, this is difficult to compute at late times, when the disc is formed and the line integral becomes inaccurate.

Note that the event horizon grows from an essentially zero size to its asymptotic value. In contrast, the apparent horizon grows from an initially non-zero size and, as it should, is always contained within the event horizon. At late times, the worldlines merge to the precision at which we can compute them.

Firstly, in the case of model D1, the differences between the equatorial and polar circumferential radii of the two trapped surfaces are very small and emerge only in the inset that offers a magnified view of the worldlines during the final stages of the collapse. This is not the case for model D4, for which the differences are much more evident and can be appreciated also in the main panel. Of course, this is what one expects given that the ratio of these two quantities depends on $a/M$ and is $\sim 1$ for a slowly rotating black hole.
Secondly, the worldlines of the stellar equatorial circumferential radius are very different in the two cases. In the slowly rotating model D1, in particular, the star collapses smoothly and the worldline always has negative slope, thus reaching progressively smaller radii as the evolution proceeds (cf. left panel of Fig. 4). By time $t \approx 0.59$ ms, the stellar equatorial circumferential radius has shrunk below the corresponding value of the event horizon. In the case of the rapidly rotating model D4, on the other hand, this is no longer true and after an initial phase which is similar to the one described for D1, the worldline does not reach smaller radii. Rather, the stellar surface slows its inward motion and, at around $t \sim 0.6$ ms, the stellar equatorial circumferential radius does not vary appreciably. Indeed, the right panel of Fig. 4 shows that at this stage the stellar surface moves to slightly larger radii. This behaviour marks the phase in which a flattened configuration has been produced and the material at the outer edge of the disc experiences a stall. As the collapse proceeds, however, also this material will not be able to sustain its orbital motion and, after $t \sim 0.7$ ms, the worldline moves to smaller radii again. By a time $t \approx 0.9$ ms, the stellar equatorial circumferential radius has shrunk below the corresponding value of the event horizon.

More detailed results as well as extensive tests of our numerical code will appear in [45].

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References

[1] Nakamura T, Oohara K and Kojima Y 1987 Prog. Theor. Phys. Suppl. 90 218
[2] Shibata M and Nakamura T 1995 Phys. Rev. D 52 5428
[3] Baumgarte T W and Shapiro S L 1999 Phys. Rev. D 59 024007
[4] Shibata M, Baumgarte T W and Shapiro S L 2000 Phys. Rev. D 61 044012
[5] Shibata M Phys. Rev. D 2003 67 024033
[6] Alcubierre M, Brügmann B, Dramlitsch T, Font J A, Papadopoulos P, Seidel E, Stergioulas N and Takahashi R 2000 Phys. Rev. D 62 62044034
[7] Font J A, Goodale T, Iyer S, Miller M, Rezzolla L, Seidel E, Stergioulas N, Suen W M and Tobias M 2002 Phys. Rev. D 65 084024
[8] Yoneda G and Shinkai H 2002 Phys. Rev. D 66 124003
[9] Kawamura M, Oohara K and Nakamura T 2003 Prog. Theor. Phys. submitted (Preprint astro-ph/0306481)
[10] Seidel E and Suen W M 1992 Phys. Rev. Lett. 69 1845
[11] Brandt S, Correll R, Gómez R, Huq M F, Laguna P, Lehner L, Marronetti P, Matzner R A, Neilsen D, Pullin J et al. 2000 Phys. Rev. Lett., 85 5496
[12] Alcubierre M and Brügmann B 2001 Phys. Rev. D 63 104006
[13] Kidder L E, Scheel M A, Teukolsky S A, Carlson E D and Cook G B 2000 Phys. Rev. D 62 084032
[14] Kidder L E, Scheel M A and Teukolsky S A 2001 Phys. Rev. D 64 064017
[15] Alcubierre M, Brügmann B, Pollney D, Seidel E and Takahashi R 2001 Phys. Rev. D 64 61501
[16] Yo H J, Baumgarte T W and Shapiro S L 2002 Phys. Rev. D 65 084026
[17] Laguna P and Shoemaker D 2002 Class. Quantum Grav. 19 3679
[18] Choptuik M W, Hirschmann E W, Liebling S L and Pretorius F 2003 Phys. Rev. D 68 044007
[19] Sperhake U, Smith K L, Kelly B, Laguna P and Shoemaker D 2004 Phys. Rev. D 69 024012
[20] May W M and White R H 1967 Methods in Computational Physics, vol 7 ed Alder B p 129
[21] Hawke I, Löffler F and Nerozzi A 2005 (Preprint gr-qc/0501054)
[22] Nakamura T and Sasaki M 1981 Phys. Lett. B, 106 69
[23] Bardeen J M and Piran T 1983 Phys. Reports 196 205
[24] Stark R F and Piran T 1985 Phys. Rev. Lett. 55 891
[25] Stark R F and Piran T 1986 Proceedings of the Fourth Marcell Grossman Meeting on General Relativity, ed Ruffini R Elsevier p 327
[26] Stark R F and Piran T 1987 Comp. Phys. Rep. 5 221
[27] Shibata M 2003 Astrophys. J. 595 992
[28] Duez M D, Shapiro S L and Yo H J 2004 Phys. Rev. D 68 104016
[29] Baiotti L, Hawke I, Montero P and Rezzolla L 2003 Mem. Soc. Astron. It. vol 1, ed Capuzzo-Dolcetta R p 327
[30] Font J A 2003 Living Rev. Relativity 6 4
[31] European RTN on Sources of Gravitational Waves, www.eu-network.org.
[32] Friedman J L, Ipser, J R and Sorkin R D 1988 Astrophys. J. 325 722
[33] Stergioulas N and Friedman J L 1995 Astrophys. J. 444 306
[34] Stergioulas N and Font J A 2001 Phys. Rev. Lett. 86 1148
[35] Stergioulas N 2003 Living Rev. Relativity 6 3
[36] Brandt S and Seidel E 1995 Phys. Rev. D 52 856
[37] Ashtekar A, Beetle C and Fairhurst S 2000 Class. Quantum Grav. 17 253
[38] Ashtekar A, Beetle C and Dreyer O, Fairhurst S, Krishnan B, Lewandowski J and Wisniewski J 2000 Phys. Rev. Lett. 85 3564
[39] Ashtekar A, Beetle C and Lewandowski J 2001 Phys. Rev. D 64 044016
[40] Ashtekar A and Krishnan B 2002 Phys. Rev. Lett. 89 261101
[41] Dreyer O, Krishnan B, Shoemaker D and Schnetter E 2002 Phys. Rev. D 67 024018
[42] Ashtekar A and Krishnan B 2003 Phys. Rev. D 68 104030
[43] Christodoulou D 1979 Phys. Rev. Lett. 25 1596
[44] Brandt S and Seidel E 1995 Phys. Rev. D 52 870
[45] Baiotti L, Hawke I, Montero P J, Löffler F L, Rezzolla L, Stergioulas N, Font, J A and Seidel E 2005 Phys. Rev. D in press (Preprint gr-qc/0403029)