Lectures on compact objects

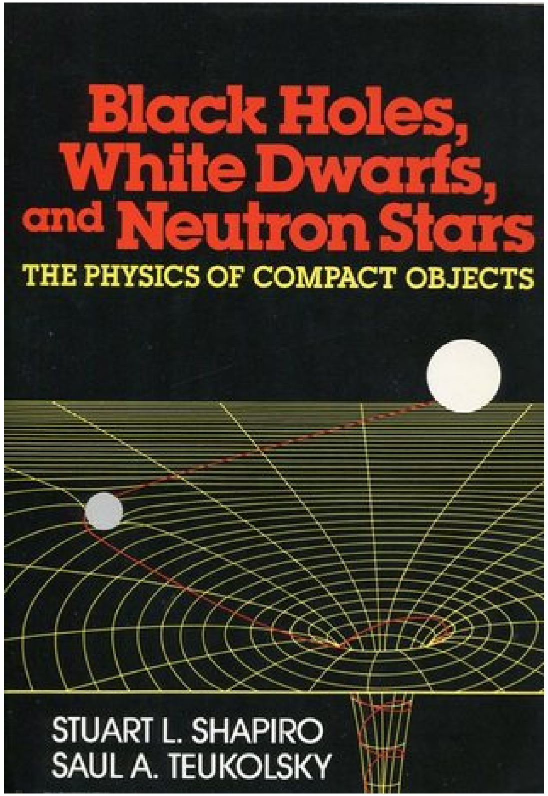
Sebastiano Bernuzzi (FSU Jena)

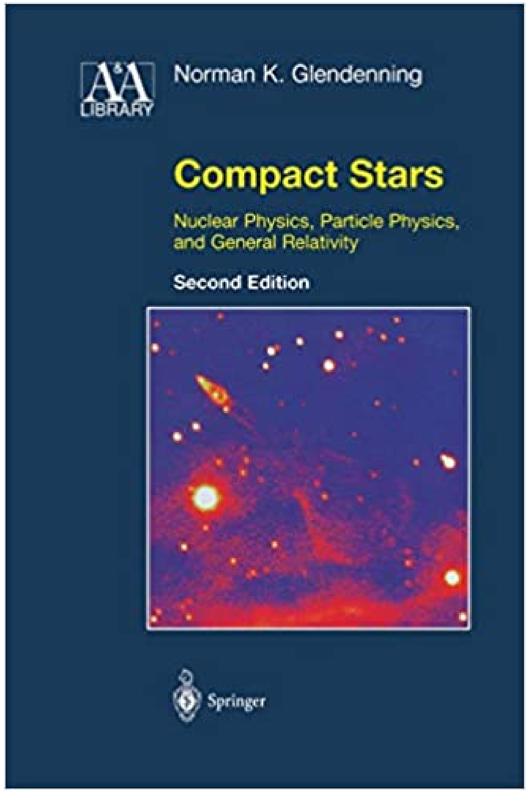
Contact: sebastiano.bernuzzi@uni-jena.de

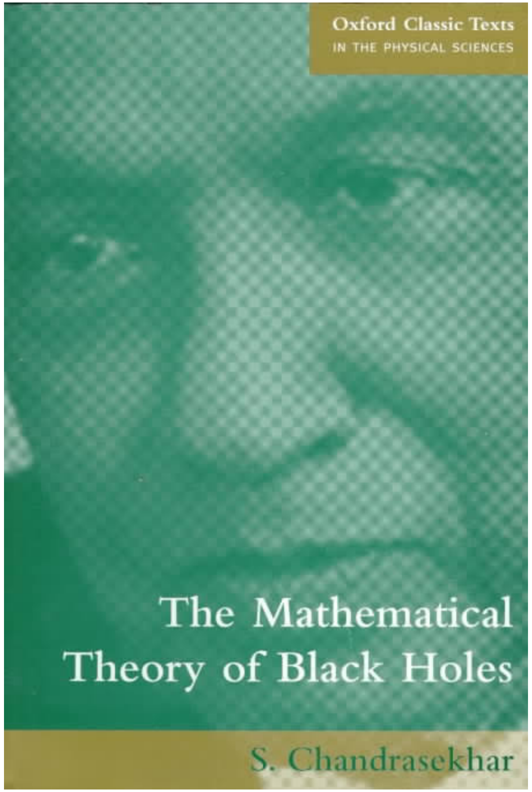
ISAPP Summer School on Gravitational Waves
June 2021

Topics

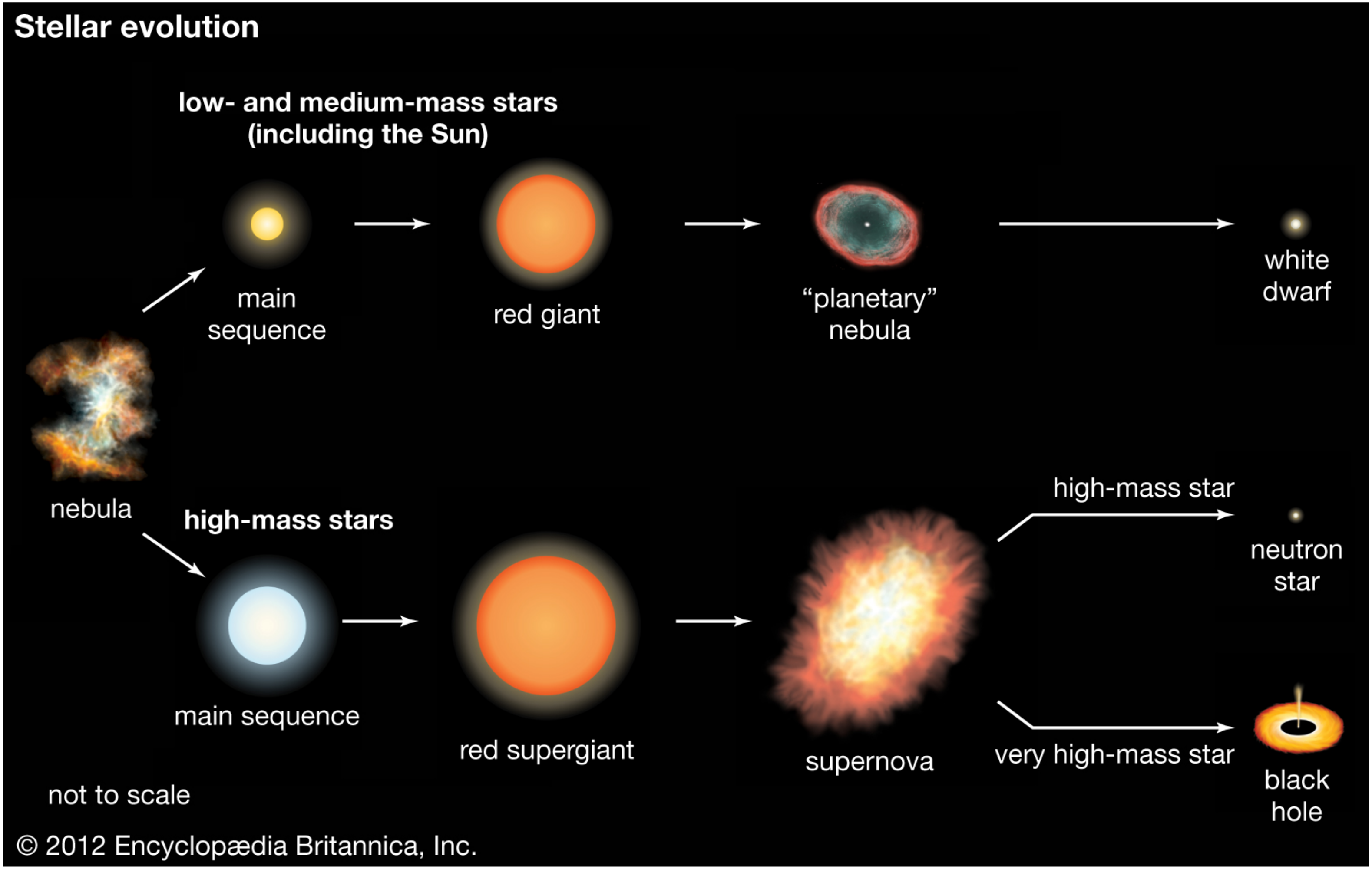
- Compact objects & compactness
- WD: Degenerate Fermi gas
- WD: Chandrasekhar mass
- NS: TOV equations & Buchdal limit
- NS: Maximum mass & Mass-radius diagram
- NS: Pulsations & stability
- NS: Equation of State
- BH: Schwarzschild solution and maximal extension
- BH: Birkhoff theorem & Schwarzschild's Orbits
- BH: the simpest GR 2-body problem
- BH: Gravitational collapse
- BH: Perturbation, Stability & Quasi-Normal-Modes
- BH (in binaries): Ringdown
- NS (in binaries): Love number, tidal polarizability & interactions


Black Holes, White Dwarfs, and Neutron Stars
THE PHYSICS OF COMPACT OBJECTS
STUART L. SHAPIRO
SAUL A. TEUKOLSKY

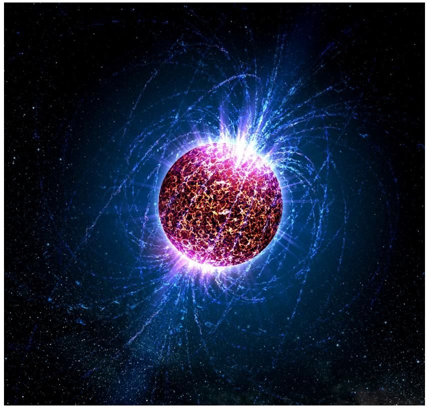

A&A LIBRARY
Norman K. Glendenning
Compact Stars
Nuclear Physics, Particle Physics, and General Relativity
Second Edition
Springer


Oxford Classic Texts
IN THE PHYSICAL SCIENCES
The Mathematical Theory of Black Holes
S. Chandrasekhar

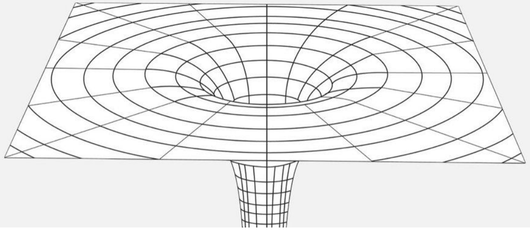
Compact objects are born when normal stars die



Compact objects

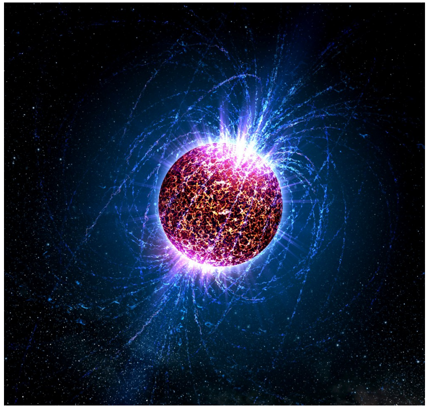


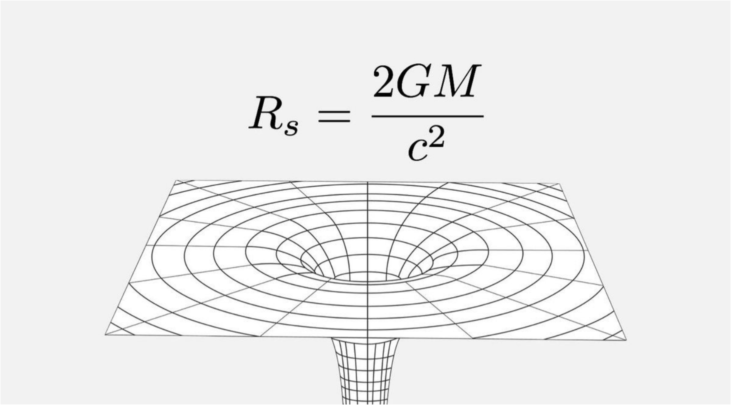
$$R_s = \frac{2GM}{c^2}$$



	White Dwarf	**Neutron Star**	**Black hole**
Progenitor Mass ($M_\odot$)	$0.1 - 8$	$8 - 25$	≥ 25
End stage of star	Red giant	Supernova	Supernova
Mass M ($M_\odot$)	$0.1 - 1.4$	$1 - 3$	≥ 3
Radius R (km)	$10^4 \sim$ Earth	10	R_{Sch}
Density ρ (kg m^{-3})	10^{9-10}	10^{18}	-
Compactness param.	10^{-4}	$0.2 - 0.4$	1
Counterbal. of self-grav.	electron deg. press.	neutron deg. press.	-

Compact objects





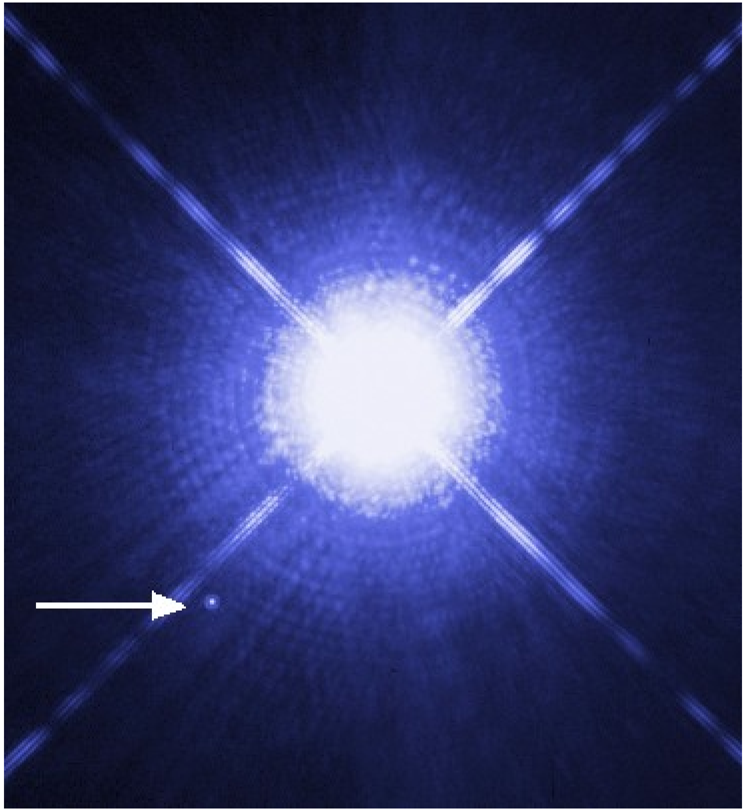
Extreme matter: Density

Extreme gravity: Surface gravity

- “cold” : do not burn fuel, no thermal pressure (degenerate matter)
- “small” : significantly smaller radii than stars with same masses

WD: Sirius B observation

- Binary system
- D ~ 8.6 light-years from Earth
- Orbital period ~ 50 yrs
- Sirius A mass ~1 M_{sun}
- Luminosity L ~ 0.06 L_{sun}
- "white" spectrum → T_{eff} ~ 25,000k
- Radius R ~ 10^{-2} R_{sun}



[HST image from wikipedia]

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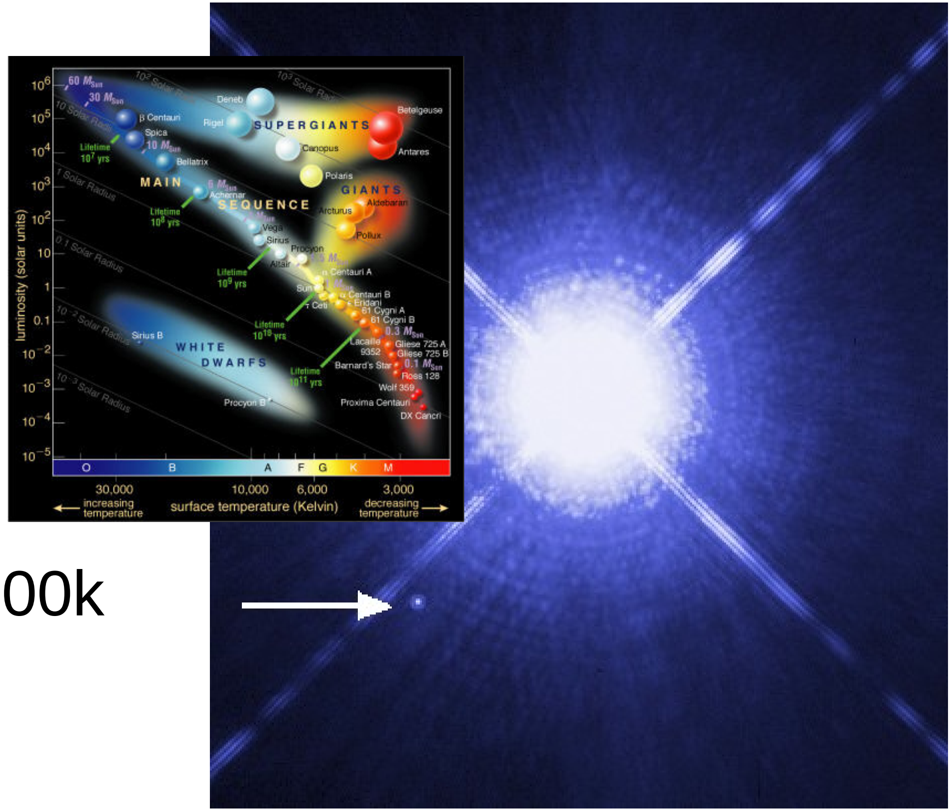
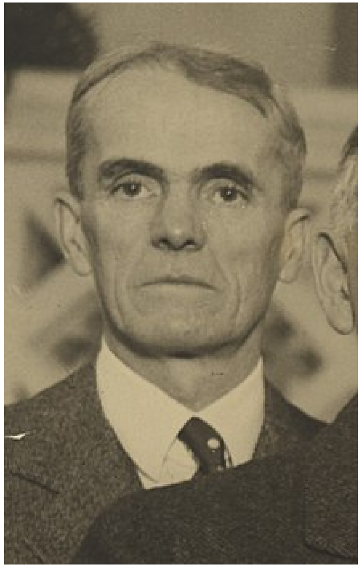


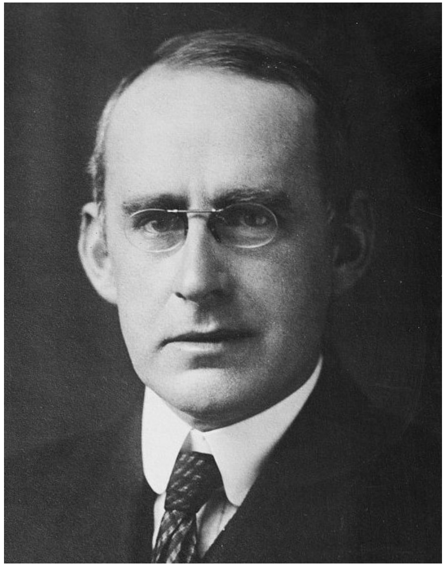
[HST image from wikipedia]



W.S.Adams (1925) Sirius B redshift measurement z~20km/s
[~89+/-19 km/s Greenstein+(1971), ~80+/-5 HST]

Prof.Adams has killed two birds with one stone; he has carried out a new test of Einstein's theory of relativity and he has confirmed our suspicion that matter 2000 times denser then platinum is not only possible, but it is actually present in the Universe

A.Eddington (1925)



It seems likely that the ordinary failure of the gas laws due to finite sizes of molecules will occur at these high densities, and I do not suppose that the white dwarfs behave like perfect gas

A.Eddington (1926)

R.H.Fowler (1926) applied the (brand-new) Fermi-Dirac statistic (Dirac 1926):
WD can be supported from gravitational collapse by **electron degeneracy pressure**

Degenerate Fermion gas EOS: definitions

Number density in phase space (# particles per unit phase space):

$$\frac{dN}{d^3x\,d^3p} = \frac{(2s+1)}{h^3} \underbrace{f(x,p)}_{\text{prob.density fun.}} \tag{1}$$

Number density:

$$n := \int \frac{dN}{d^3x\,d^3p}\, d^3p \tag{2}$$

Energy density ($E^2 = pc^2 + mc^2$) :

$$\varepsilon := \int \frac{dN}{d^3x\,d^3p}\, E d^3p \tag{3}$$

(Isotropic) Pressure:

$$P := \frac{1}{3}\int \frac{dN}{d^3x\,d^3p}\, \frac{p^2 c^2}{E}\, d^3p \tag{4}$$

At equilibrium the prob.density function is the Fermi-Dirac distribution:

$$f = \frac{1}{e^{\beta(E-\mu)+1}}\,, \quad \text{with} \quad \beta := \frac{1}{k_B T} \tag{5}$$

Degenerate Fermion gas EOS: Ideal electrons gas

Limits of the Fermi-Dirac

- high-temp/low-density ($E \ll k_B T$): $f(E) \approx e^{\beta(E-\mu)} \ll 1$
- low-temp/high-density: $f(E) = \theta(E_F - E)$ with $E_F := \mu(0)$ the **Fermi energy**

Work in the low-temp limit, compute number density, pressure & energy ($\int d^3p = 4\pi \int p^2\, dp$):

$$n_e = \frac{8\pi}{h^3} \int_0^{p_F} p^2 dp = \frac{8\pi}{3h^3} p_F^3 = \frac{8\pi}{3h^3} (m_e c)^2 x^3 \quad \text{with} \quad x := \frac{p_F}{m_e c} \tag{6}$$

$$P_e = \frac{8\pi}{3h^3} \int_0^{p_F} \frac{p^2 c^2}{(p^2 c^2 + m^2 c^4)^{1/2}} p^2 dp = \frac{8\pi}{3h^3} m_e^4 c^5 \underbrace{\int_0^x \frac{y^4}{(1-y^2)^{1/2}} dy}_{=:\phi(x)} = \frac{m_e^4 c^5}{\hbar^3} \phi(x) \tag{7}$$

with

$$\phi(x) = \frac{1}{8\pi^2} \left[x (1+x^2)^{1/2} \left(\frac{2}{3} x^2 - 1 \right) + \ln\left[x + (1+x^2)^{1/2}\right] \right] \tag{8}$$

Similarly:

$$\varepsilon_e = \frac{m_e^4 c^5}{\hbar^3} \chi(x) \tag{9}$$

$$\chi(x) = \frac{1}{8\pi} [x(1+x^2)^{1/2} (1+2x^2) - \ln[x + (1+x^2)^{1/2}]] \tag{10}$$

Degenerate electron gas EOS: mass density & polytropic limits

In a system with electrons and ions, the charge of each electron is neutralized by a proton which is turn is accopanied by one or more neutrons in the nucleous. The rest-mass density of the gas is dominated by the mass of the nucleons since $m_e \ll m_N$; indicating as A/Z the number of nucleons per electron,

$$\rho_0 = \sum_{\mathrm{species},s} n_s m_s \approx n_e \frac{A}{Z} m_N \tag{11}$$

Total mass density:

$$\rho = \rho_0 + \varepsilon_e / c^2 \approx \rho_0 \tag{12}$$

The EOS is parametrically given by Eq.(6),(7),(9) and (11) as functions of x .

For example, by combining Eq.(11) with Eq.(6) the Fermi parameter (momentum) is

$$x = \frac{h}{m_e c}\left(\frac{3}{8\pi}\frac{Z}{A}\frac{\rho}{m_N}\right)^{1/3} \tag{13}$$

and the other quantities can be expressed as $P_e(\rho)$ and $\varepsilon_e(\rho)$.

Note: $A/Z \sim 2$ for He, C, O.

Define the *critical density*

$$\rho_c := \frac{1}{3\pi^2}\frac{A}{Z}\Big(\frac{m_e c}{\hbar}\Big)^3 \sim \frac{A}{Z} 10^6 \, g/\mathrm{cm}^3 \tag{14}$$

- Low density limit, $\rho \ll \rho_0$: non-relativistic electrons $cp_F \ll m_e c^2$ or $x \ll 1$

$$\phi\,(x) \sim \frac{1}{15\pi^2} x^5 \qquad \chi\,(x) \sim \frac{1}{3\pi^2} x^3 \tag{15}$$

- High density limit, $\rho \gg \rho_0$: relativistic electrons $x \gg 1$

$$\phi\,(x) \sim \frac{1}{12\pi^2} x^4 \qquad \chi(x) \sim \frac{1}{4\pi^2} x^4 \tag{16}$$

In these limits the EOS can be written in the form

$$P_e = K \rho_0^\Gamma \qquad \varepsilon_e = \kappa \rho_0 \tag{17}$$

with (different) constants K, Γ and κ:

- $\rho \ll \rho_c$: $\Gamma = 5/3$ $\qquad K = \big(\frac{3}{8\pi}\big)^{2/3} \frac{h^2}{5\,m_e} \Big(\frac{Z}{A m_N}\Big)^{5/3} \sim 10^{13} \Big(\frac{Z}{A}\Big)^{5/3} [\mathrm{CGS}]$ $\qquad \kappa = \frac{3}{2}$
- $\rho \gg \rho_c$: $\Gamma = 4/3$ $\qquad K = \big(\frac{3}{\pi}\big)^{1/3} \frac{c\,h}{8} \Big(\frac{Z}{A m_N}\Big)^{4/3} \sim 10^{15} \Big(\frac{Z}{A}\Big)^{4/3} [\mathrm{CGS}]$ $\qquad \kappa = 3$

Degenerate neutron gas EOS

The above derivation is valid for any gas made of non-interacting fermions, by appropriately substituting the value of the particle mass. For example, a gas of non-interacting neutrons ($m_n/m_e \sim 10^3$) has a critical density of

$$\rho_c := \frac{1}{3\pi^2}\left(\frac{m_n c}{\hbar}\right)^3 \sim 6\times 10^{15}\, g/\mathrm{cm}^3 \tag{18}$$

and it is described by a polytropic EOS with

- $\rho \ll \rho_c : \Gamma = 5/3 \qquad K = \frac{3^{2/3}\pi^{4/3}}{5}\frac{\hbar^2}{m_n^{8/3}} \sim 10^9\,[\mathrm{CGS}]$
- $\rho \gg \rho_c : \Gamma = 4/3 \qquad K = \frac{3^{1/3}\pi^{2/3}}{4}\frac{\hbar c}{m_n^{4/3}} \sim 10^{15}\,[\mathrm{CGS}]$

In this case the mass density $\rho = \rho_0 + \varepsilon_n/c^2 \approx \varepsilon_n/c^2$ is dominated by the neutron's energy density in the relativistic limit.

Fate of a star with “no fuel”

Without nuclear burning a star decreases its total energy E because of radiation emission from the surface. The energy source in this phase is the gravitational energy $W \sim -GM/R$. The virial theorem implies that the star contracts

$$\Delta E < 0 \Rightarrow \Delta R < 0 \tag{19}$$

For a Maxwell-Boltzmann ideal gas $P \propto k_B T$, the gravitational energy is proportional to the average temperature of the star:

$$-W = 4\pi \int_0^R \frac{Gm(r)\rho(r)}{r} r^2 dr = 4\pi \int_0^R \frac{dP}{dr} r^3 dr = -12\pi \int_0^R P r^2 dr \propto \bar{T} \tag{20}$$

(The second equality uses the hydrostatic equilibrium equation, and the third integrates by part). This implies that, as the radius decreases, the temperature increases as

$$\bar{T} \sim \frac{M}{R} \tag{21}$$

At the same time the average density increases as

$$\bar{\rho} \sim \frac{M}{R^3} \tag{22}$$

Under these conditions a gas of $M \sim M_\odot$ becomes degenerate and the star cannot contract for ever: Pauli’s exclusion principle becomes relevant and the Fermi-Diract statistics must be applied.

Indeed, the typical momentum of the electrons in a Maxwell-Boltzmann gas is estimated as

$$\Delta p_e \sim (k_B \bar{T})^{1/2} \sim \left(\frac{M}{R}\right)^{1/2} \tag{23}$$

while their separation is

$$\Delta q_e \sim \left(\frac{1}{\bar{\rho}}\right)^{1/3} \sim \frac{R}{M^{1/3}} \tag{24}$$

The volume in the phase space is thus

$$(\Delta p_e \Delta q_e)^3 \sim \left(\frac{M^{1/2}}{R^{1/2}} \frac{R}{M^{1/3}}\right)^3 \sim (R^{1/2} M^{1/6})^3 \tag{25}$$

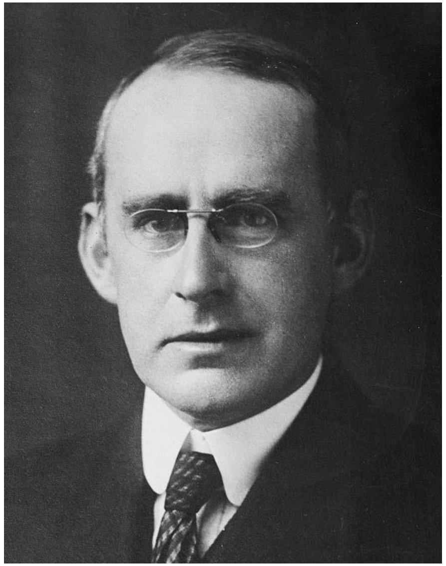
and with all the factors ...

$$(\Delta p_e \Delta q_e)^3 \sim 40 \left[10^{26} \left(\frac{M}{M_\odot}\right)^{1/6} \left(\frac{R}{R_\odot}\right)^{1/2} [\mathrm{CGS}]\right]^3 \sim h^3 \left(\frac{M}{M_\odot}\right)^{1/2} \left(\frac{R}{3 \times 10^{-2} R_\odot}\right)^{3/2} \tag{26}$$

Hence, the electron degeneracy pressure starts to support a star of a solar mass that contracts to $\sim 3 \times 10^{-2} R_\odot$

Exercises

- Derive yourself all the discussed results about the degenerate electron gas
- Use the virial theorem to argue that a star without fuel must contract
- Estimate the gravitational redshift of a photon emitted from the source of a WD



Chandrasekhar … has shown that a star of a mass greater than a certain limit remains a perfect gas …

The star has to go on radiating and radiating and contracting and contracting until, I suppose, it gets down to a few km radius, when gravity becomes strong enough to hold in the radiation, and the star can at last find peace. ... I think there should be a law of Nature to prevent a star from behaving in this absurd way!

A.Eddington (1935)

It is clear from this statement that Eddington fully realized, already in 1935, that given the existence of an upper limit to the mass of degenerate configurations, one must contemplate the possibility of gravitational collapse leading to the formation of what we now call black holes. But he was unwilling to accept a conclusion that he so presciently drew.

S.Chandrasekhar (1982)

Landau's argument (1932)

Consider N fermions in a star of radius R, so that the number density is $n = N/R^3$, $\Delta q \sim (1/n)^3$ and, from the Heisemberg uncertainty principle, $\Delta p \sim \hbar n^{1/3}$. The Fermi energy of the gas in the relativistic regime is

$$E_F \sim \hbar c n^{1/3} \sim \frac{\hbar c N^{1/3}}{R} \tag{1}$$

The mass of the star is dominated by the baryons $M \sim N m_B$ and the gravitational energy per Fermion is

$$W \sim -\frac{G M m_B}{R} = -\frac{G m_B^2 N}{R} \tag{2}$$

Define the equilibrium state as the minimum of the total energy

$$E = E_F + W = \frac{\hbar c N^{1/3}}{R} - \frac{G m_B N}{R} \tag{3}$$

There are two cases

i. N small, $E > 0$: The energy can be decreased by increasing R (decreasing E_F). When W dominates, E becomes negative and then tends to zero for large R. Given any value of R a minimum can be found: a stable equilibrium always exists.

ii. N large, $E<0$: The energy can be decreased without bound by decreasing R. No equilibrium can exists and gravitational collapse happens.

The maximum baryon number (and thus mass) that guarantees equilibrium is then determined by setting $E=0$:

$$N_{\max}\sim\left(\frac{\hbar c}{G m_B}\right)^{3/2}\sim 2\times 10^{57} \quad \Rightarrow M_{\max}\sim N_{\max}\, m_B\sim 1.5\, M_{\odot} \tag{4}$$

$\Rightarrow$ The order of magnitude of the maximum mass is determined by fundamental constants.

The equilibrium radius associated to $M_{\max}$ can be estimated by requiring that the instability happens as matter becomes relativistic

$$E_F\gtrsim m c^2 \tag{5}$$

with m the particle mass. Using Eq.(4) in Eq.(1) one gets

$$R\lesssim\frac{\hbar}{mc}\left(\frac{\hbar c}{G m_B}\right)^{1/2} \tag{6}$$

This gives $R\sim 5\times 10^8\,\mathrm{cm}$ for electrons ($m=m_e$) and $R\sim 3\times 10^5\,\mathrm{cm}$ for neutrons ($m=m_n$).

Chandrasekhar's maximum mass

Another simple argument for the maximum mass.

Assume a polytropic EOS $P \propto \rho^\Gamma$. The pressure gradient scales as

$$P \sim \frac{M^\Gamma}{R^{3\gamma}} \quad \Rightarrow \quad \frac{dP}{dr} \sim M^\Gamma R^{-3\Gamma-1} \tag{7}$$

The gravitational force scales as

$$-\frac{G m \rho}{r^2} \sim M^2 R^{-5} \tag{8}$$

Impose equilibrium

$$\frac{dP}{dr} = -\frac{G M \rho}{r^2} \quad \Rightarrow \quad M^\Gamma R^{-3\Gamma-1} \sim M^2 R^{-5} \tag{9}$$

and

$$R \sim M^{\frac{\Gamma-2}{3\Gamma-4}} \tag{10}$$

This equation gives the *correct* scaling of the radius with mass for Newtonian polytropes and can be proved using the Lane-Emden equation.

For $\Gamma \to 4/3$ (all electrons are relativistic) the radius tends to zero and no star exists. In fact, using the Lane-Emden equation it is possible to show that in this limit M becomes *independent* on the central density and attains the value

$$M_{\rm Ch} = 5.74 \left(\frac{A}{Z}\right)^2 M_\odot \sim 1.435\, M_\odot \; (\text{for } A/Z = 2) \tag{11}$$

In practice, however, the maximum mass is reached at finite density, and stable WD sequences end as a consequence of a compositional change in the star. The main responsibile for this is *neutronization*: at sufficiently high densities $\rho \sim 10^9\, g/\text{cm}^3$ inverse β-decay reactions

$$p + e^- \longrightarrow n + \nu_e \tag{12}$$

take place and produce nuclei rich of neutrons, e.g.

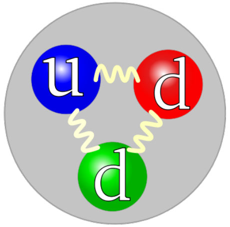
$$^{56}_{26}\text{Fe} \longrightarrow {}^{56}_{25}\text{Mn} \longrightarrow {}^{56}_{24}\text{Cr} \tag{13}$$

As a consequence, the EOS softens (less pressure support) and the adiabatic index decreases $\Gamma < 4/3$ leading to a collapse. Neutronization is the same process that in the core of contracting massive stars (at higher densities) leads to the formation of stable NSs.

Note that WD at densities $\rho \sim 10^{10}\, g/\text{cm}^3$ are also subject to a gravitational instability.

J.Chadwick (1932) discovers the neutron

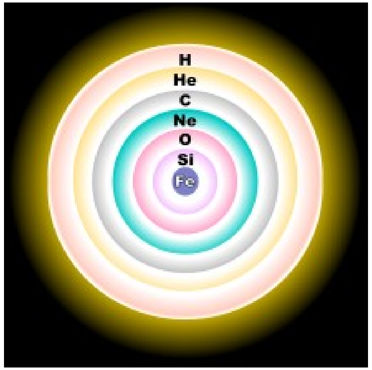
m = 1,674 927 351(74) × 10^{-24} g



With all reserve we advance the view that supernovae represent the transition from ordinary stars into neutron stars, *which in their final stages consist of closely packed neutrons*

W.Baade & F.Zwicky (1933)

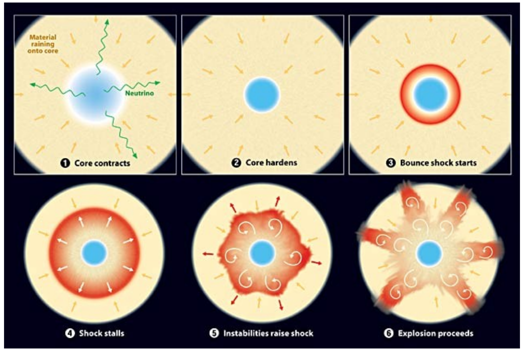
Core collapse of massive stars



$$^{56}\mathrm{Ni} \longrightarrow {}^{56}\mathrm{Fe} + 2e^{+} + 2\nu_e$$

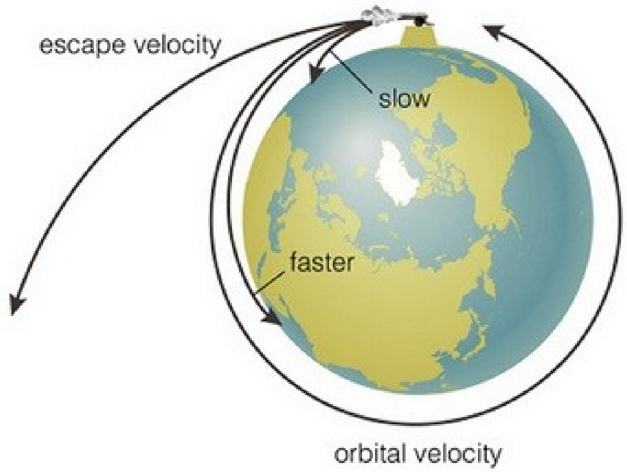
$$^{A}\mathrm{Z} + e^{-} \longrightarrow {}^{A}(\mathrm{Z\text{-}1}) + \nu_e$$

$$^{56}\mathrm{Fe} + \gamma \longrightarrow 13\mathrm{He} + 4n$$



- Core has $\Gamma < 4/3$ => collapse
- $8 < M/M_{Sun} < 30$ => neutron degeneracy pressure and repulsive nucleon-nucleon interaction stabilize the core ($\Gamma > 4/3$, $\rho \sim 10^{14}$ g/cm^3)
- Core bounce and supernova explosion

$$\Delta W \sim -G\left(\frac{M_{\mathrm{core}}}{1.4\,M_\odot}\right)^2\left(\frac{R_{\mathrm{core}}}{10^9\,\mathrm{cm}} - \frac{R_{\mathrm{NS}}}{10^6\,\mathrm{cm}}\right)^{-1} \sim -G\left(\frac{M_{\mathrm{core}}}{1.4\,M_\odot}\right)^2\left(\frac{R_{\mathrm{NS}}}{10^6\,\mathrm{cm}}\right)^{-1} \sim 10^{53}\,\mathrm{erg}$$

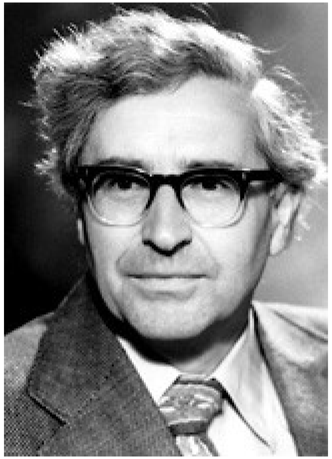


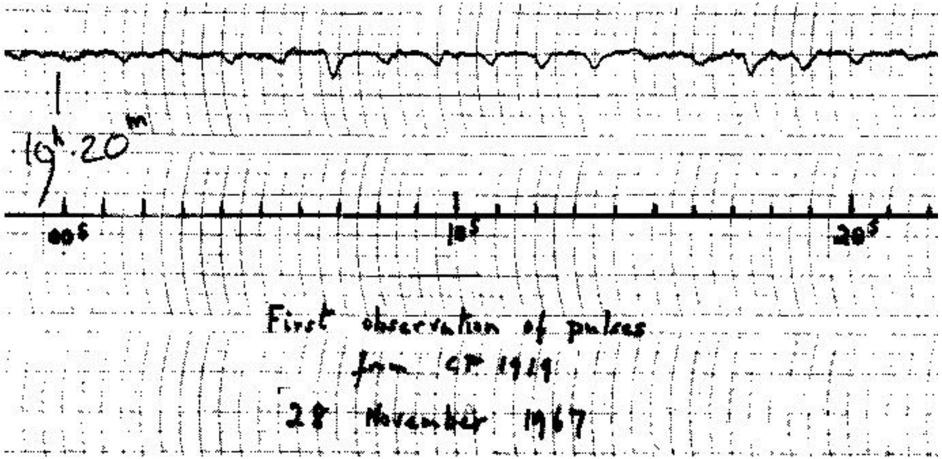
Surface grav. field = 10^{11} x Earth
Escape velocity ~ 1/3 c

$$v_{\rm esc} = \sqrt{\frac{2\,G\,M}{R}}$$

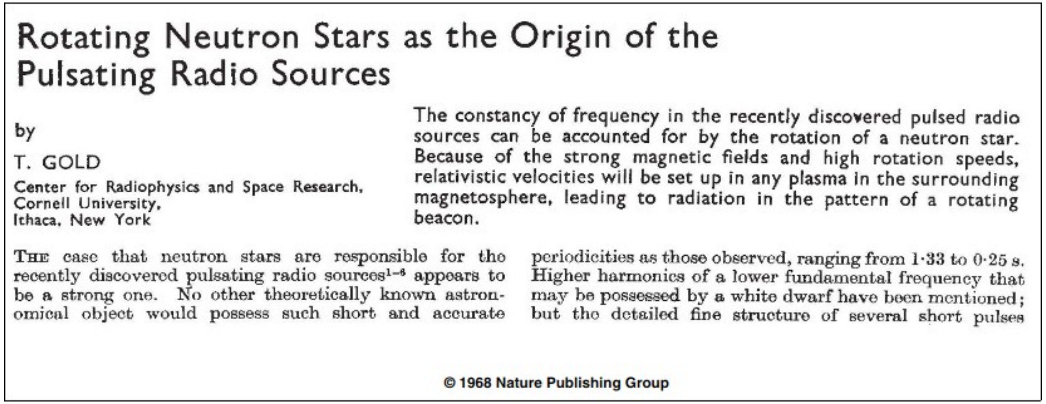
Must be described with general relativity!

Pulsars





Bell / Hewish 1968 (Ryle & Hewish Nobel prize 1974)



Rotating Neutron Stars as the Origin of the Pulsating Radio Sources

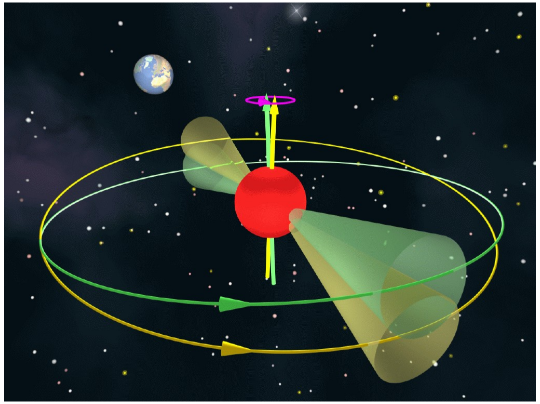
by

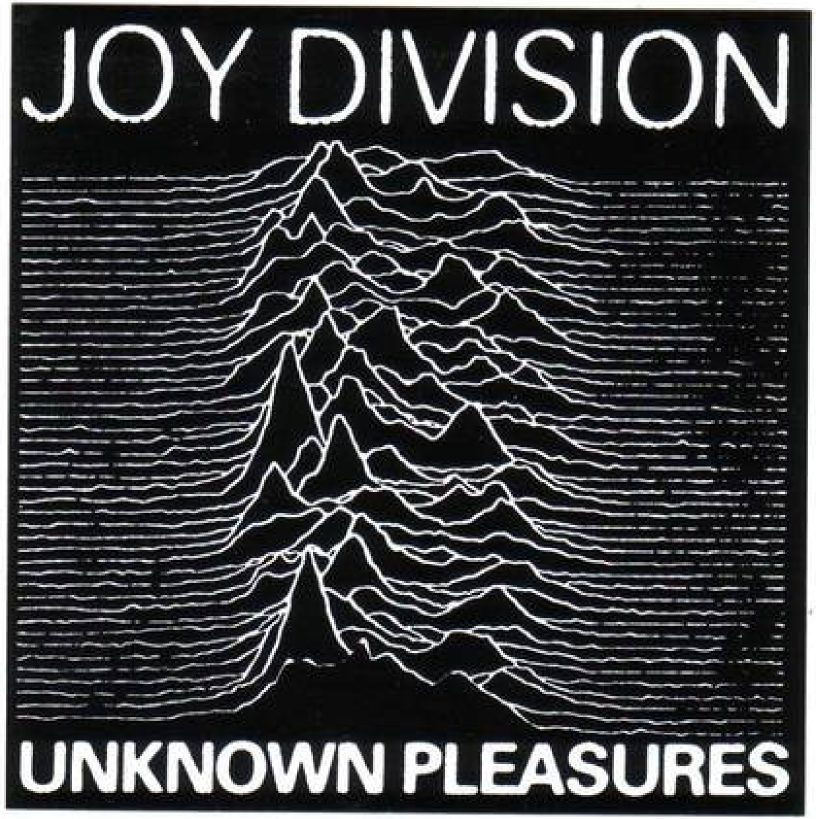
T. GOLD

Center for Radiophysics and Space Research,
Cornell University,
Ithaca, New York

The constancy of frequency in the recently discovered pulsed radio sources can be accounted for by the rotation of a neutron star. Because of the strong magnetic fields and high rotation speeds, relativistic velocities will be set up in any plasma in the surrounding magnetosphere, leading to radiation in the pattern of a rotating beacon.

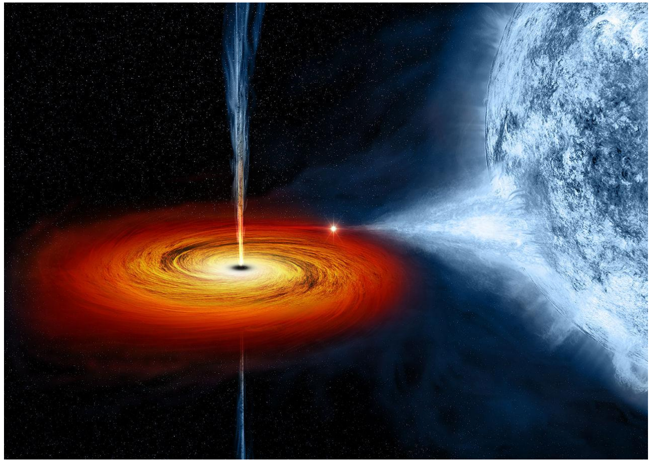
THE case that neutron stars are responsible for the recently discovered pulsating radio sources[1–6] appears to be a strong one. No other theoretically known astronomical object would possess such short and accurate periodicities as those observed, ranging from 1·33 to 0·25 s. Higher harmonics of a lower fundamental frequency that may be possessed by a white dwarf have been mentioned; but the detailed fine structure of several short pulses

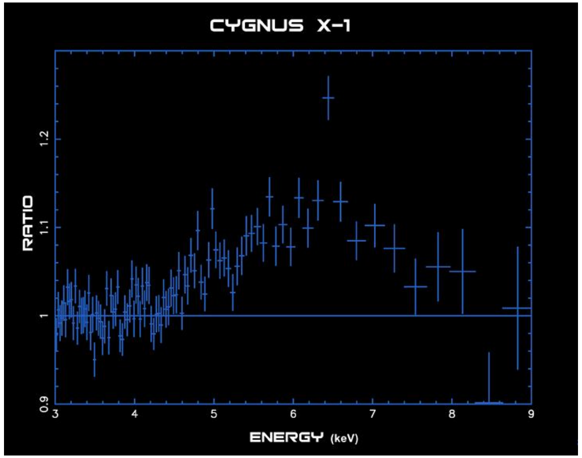



JOY DIVISION
UNKNOWN PLEASURES

X-ray astronomy

- Hot gases at T ~ 1,000,000K emit X-ray
- 1962 Scorpius X-1
  - Strongest X-ray source together the Sun.
  - Low-Mass-X-ray binary
  - $1.4M_{Sun}$ NS + 0.42 star
- 1964 Cygnus X-1
  - High-Mass-X-ray binary
  - $14.8M_{Sun}$ BH + $20\text{-}40M_{Sun}$ supergiant star





R.Giacconi (Nobel Prize 2002)

Crab & Vela pulsars

Pulsating Radio Sources near the Crab Nebula

Abstract. *Two new pulsating radio sources, designated NP 0527 and NP 0532, were found near the Crab Nebula and could be coincident with it. Both sources are sporadic, and no periodicities are evident. The pulse dispersions indicate that 1.58 ± 0.03 and 1.74 ± 0.02 × 10^{20} electrons per square centimeter lie in the direction of NP 0527 and NP 0532, respectively.*

VOLUME 22, NUMBER 7 PHYSICAL REVIEW LETTERS 17 FEBRUARY 1969

CRAB NEBULA PULSAR NPO527

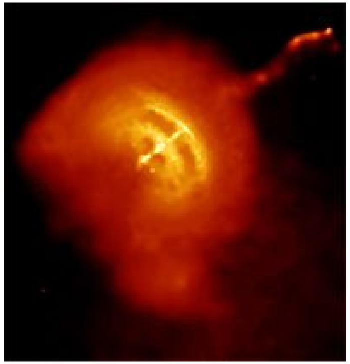
Edward C. Reifenstein, III, William D. Brundage, and David H. Staelin*
National Radio Astronomy Observatory,† Green Bank, West Virginia
(Received 30 December 1968)

Position measurements of the Crab nebula pulsar NPO527 made with a split-beam antenna yield the position (1950) $\alpha = 05^h 26^m 10^s \pm 40^s$, $\delta = 22°0' \pm 30'$. Thus NPO527 lies 1.2° from the pulsar NPO532, which is located at the nebula. The proximity and similar dispersion of these two pulsars suggest that they may have had a common origin in the supernova explosion of A.D. 1054 and hence that NPO527 is moving with a velocity of ~0.15c, a hypothesis which could be tested directly by proper motion measurements.

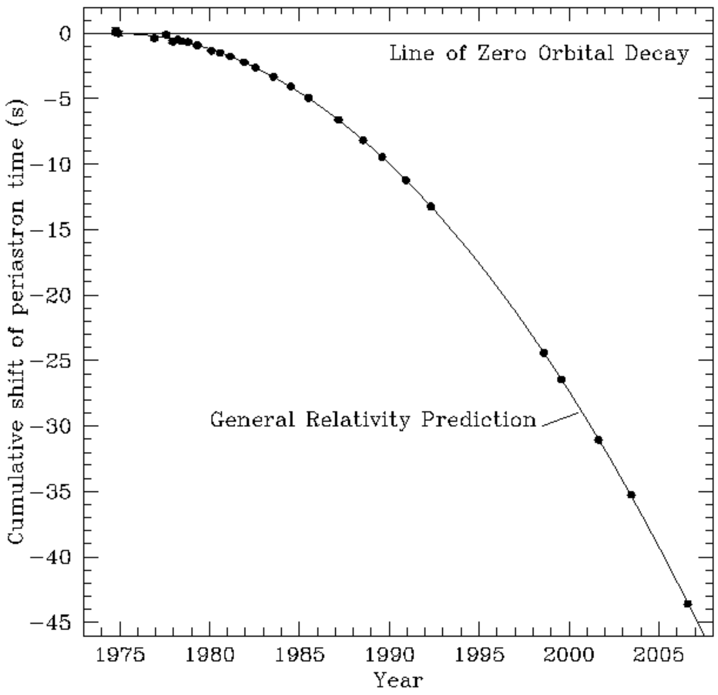
A Pulsar Supernova Association?

by
M. I. LARGE
A. E. VAUGHAN
B. Y. MILLS
Cornell-Sydney University Astronomy Centre,
School of Physics,
University of Sydney

A pulsar with a very short period (0·089 s) has been discovered at the position of a suspected supernova remnant, raising several interesting consequences.



Neutron stars in binary systems



PSR B1913+16

[Weisberg&Taylor 2004]

1993 Nobel prize: Hulse & Taylor

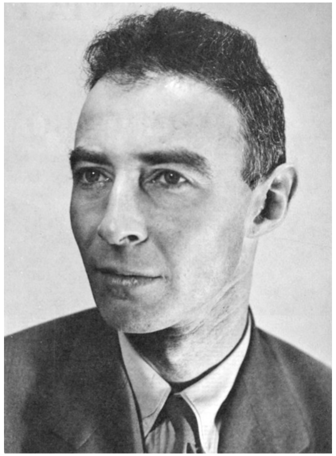
FEBRUARY 15, 1939 PHYSICAL REVIEW VOLUME 55

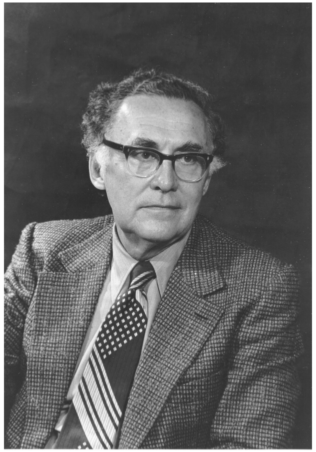
On Massive Neutron Cores

J. R. Oppenheimer and G. M. Volkoff
Department of Physics, University of California, Berkeley, California

(Received January 3, 1939)

It has been suggested that, when the pressure within stellar matter becomes high enough, a new phase consisting of neutrons will be formed. In this paper we study the gravitational equilibrium of masses of neutrons, using the equation of state for a cold Fermi gas, and general relativity. For masses under $\frac{1}{3}\odot$ only one equilibrium solution exists, which is approximately described by the nonrelativistic Fermi equation of state and Newtonian gravitational theory. For masses $\frac{1}{3}\odot < m < \frac{3}{4}\odot$ two solutions exist, one stable and quasi-Newtonian, one more condensed, and unstable. For masses greater than $\frac{3}{4}\odot$ there are no static equilibrium solutions. These results are qualitatively confirmed by comparison with suitably chosen special cases of the analytic solutions recently discovered by Tolman. A discussion of the probable effect of deviations from the Fermi equation of state suggests that actual stellar matter after the exhaustion of thermonuclear sources of energy will, if massive enough, contract indefinitely, although more and more slowly, never reaching true equilibrium.





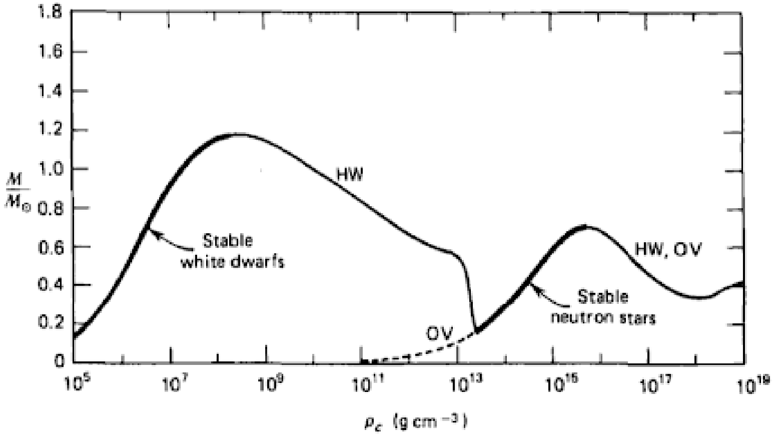


Figure 9.1 Gravitational mass versus central density for the HW (1958) and OV (1939) equations of state. The stable white dwarf and neutron star branches of the HW curve are designated by a *heavy solid* line.

Tolmann-Oppenheimer-Volkoff equations (1939)

Metric (Units: $c = G = 1$):

$$g = -e^{2\phi(r)}\,dt^2 + e^{2\lambda(r)}dr^2 + r^2 d\Omega^2 \qquad e^{2\lambda(r)} =: \left(1 - \frac{2m(r)}{r}\right)^{-1} > 0 \tag{1}$$

Perfect fluid matter (Note $\rho = \varepsilon / c^2 = \varepsilon$): $T_{\alpha\beta} = (\rho + P)u_\alpha u_\beta + P g_{\alpha\beta}$

Structure equations valid in the star *interior*

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

$$\frac{dP}{dr} = -\frac{\rho m}{r^2}\left(1 + \frac{P}{\rho}\right)\left(1 + \frac{4\pi r^3 P}{m}\right)\left(1 - \frac{2m}{r}\right)^{-1} \tag{3}$$

$$\frac{d\phi}{dr} = -\frac{1}{\rho}\left(1 + \frac{P}{\rho}\right)^{-1}\frac{dP}{dr} \tag{4}$$

augmented by a (barotropic) EOS

$$P = P(\rho) \tag{5}$$

Eq.(2) is formally equivalent to the mass-density equation in Newton gravity. The quantity

$$M = 4\pi \int_0^R \rho(r)\, r^2\, dr \tag{6}$$

is the *gravitational mass* of the NS (and matches the mass of the exterior Schwarzschild solution). However the total energy of the spacetime is given by

$$E = \int_\Sigma \sqrt{\gamma}\, T_{\alpha\beta}\, t^\alpha n^\beta d^3x = 4\pi \int_0^R \rho(r)\, e^{\lambda(r)}\, r^2\, dr > M \quad \text{because} \quad e^\lambda > 1\; (m > 0) \tag{7}$$

the difference $E - M =: E_b$ is interpreted as the binding energy of the star.

The Newtonian limit of Eq.(3) and Eq.(4) is given by $P \ll \rho$ and $2m/r \ll 1$.

Eq.(3) describes the hydrostatic equilibrium. The R.H.S. contains three correction factors with respect to the relative Newtonian equation and they all enhance the "active/effective mass" between radial shells (Pressure contributes to energy in GR!)

Eq.(4) in the Newtonian limit returns the Newtonian gravitational potential since $e^{2\phi} \approx 1 + 2\phi$, and Eq.(1) can be compared to the weak field metric.

How to obtain a solution?

- Prescribe the EOS and the central pressure (or energy density) value at $r=0$
- Evaluating the proper radius $\int_0^r e^{\lambda(x)}dx \approx e^{\lambda(0)}r$ for small radii and requiring smoothness at $r=0$ and local flatness implies that $e^{\lambda(0)}=1$ and $m(0)=0$
- Integrate the structure equations using the EOS up to the surface: $P(R)=0$
- Match $\phi(r=R)$ to the Schwarzschild vacuum exterior solution (Birkhoff theorem)

If $\rho>0$ and $dP/d\rho>0$, the TOV equations (2-3) guarantee that the pressure profile drecreases integrating outwards to the surface. Requiring $dP/d\rho>0$ guarantees stability: a fluctuation in $\rho(r=r_0)$ leading to a decrease of P would otherwise cause the nearby fluid element to move to r_0 leading to a further increase in $\rho(r_0)$ and a runaway instability.

The definition of $m(r)$ implies that $2m(r)/r<1$ for all r. Evaluating the bound at $r=R$ one obtains that *any static, spherically symmetric star must have a radius greater than its Schwarzschild radius :*

$$R>2M \tag{8}$$

There is no Newtonian analogue of this result because $2\mathrm{GM}/(c^2R)\to 0$ for $c\to\infty$ and Eq.(8) is trivially $1>0$.

Schwarzschild solution (constant density star)

Assuming a constant density star $\rho(r<R)=\rho_0$ (and $\rho=0$ for $r\geq R$), Eq.(2) can be immediately integrated to $m(r)=4/3\pi r^3\rho_0$ and it holds also in the Newtonian limit.

Substituting $m(r)$ in the pressure equation and integrating, one obtains in the *Newtonian limit*

$$P(r)=\frac{2}{3}\pi\rho_0^2\,(R^2-r^2) \tag{9}$$

that implies the central pressure is

$$P(0)=\frac{2}{3}\pi\rho_0^2R^2=\left(\frac{\pi}{6}\right)^{1/3}M^{2/3}\rho_0^{4/3} \tag{10}$$

$\Rightarrow$ The Newtonian central pressure is always finite (for any M, R).

This is not true in GR! The pressure solution in GR is

$$P(r)=\rho_0\frac{\left(1-\frac{2M}{R}\right)^{1/2}-\left(1-\frac{2M}{R^3}r^2\right)^{1/2}}{\left(1-\frac{2M}{R^3}r^2\right)^{1/2}-3\left(1-\frac{2M}{R}\right)^{1/2}} \tag{11}$$

which gives the central pressure

$$P(0) = \rho_0 \frac{1 - \left(1 - \frac{2M}{R}\right)^{1/2}}{3\left(1 - \frac{2M}{R}\right)^{1/2} - 1} \tag{12}$$

The pressure is infinite if the denominator is zero, which lead to the *Buchdal inequality*

$$R > \frac{9}{4} M \tag{13}$$

⇒ No star can esists in GR with a smaller radius. Note this bound is stronger than Eq.(8).

Buchdal's limit

The Buchdal inequality is actually a general result (for any EOS!).

Under the hypotheses $\rho \geq 0$ and $d\rho/dr \leq 0$ and using the TOV equations it is possible to establish the bound (see e.g. Wald's book):

$$\frac{m(r)}{r} < \frac{2}{9}[1 - 6\pi r^2 P(r) + (1 + 6\pi r^2 P(r))^{1/2}] \tag{14}$$

which gives immediately Eq.(13) if evaluated at $r = R$ ($P(r) = 0$).

Maximum NS mass

Analogously to WD, one expects the existance of a maximum mass also for NS. This expectation is correct, and the precise value of $M_{\max}$ depends on the NS EOS at high densities. The latter is not precisely known but under rather general hypothesis it is possible to estimate upper bounds.

Assume the EOS is known below a certain density ρ_0 attaeined at $r=r_0$ and define the *core* of the star as the $r>r_0$ region such that $\rho(r)>\rho_0$ and the envelop as the $r<r_0$ such that $\rho(r)<\rho_0$. From the TOV equations, the mass of the core $m_0=m(r_0)$ is certainly larger than

$$m_0 \geq \frac{4}{3}\pi r^3 \rho_0 \tag{15}$$

but the ratio $m_{0/}r_0$ is also bound by Eq.(14) evaluated at $r=r_0$. Since Eq.(14) is a decreasing function of P a simpler (but weaker) inequality is given by taling $P_0<0$ (Buchdal-like inequality)

$$m_0 \leq \frac{4}{9} r_0 \tag{16}$$

The two inequalities define the "existance region" of the NS in the plane (r_0, m_0) and the upper bound

$$m_0 < \sqrt{\frac{16}{243\pi\rho_0}} \tag{17}$$

Experimentally, the EOS of matter at high densities is reasonably known up to the density of atomic nuclei, $\rho_{\rm nuc} \simeq 4.6 \times 10^{14}\, g/{\rm cm}^3$. Using the "known" EOS *up to these densities* to integrate the TOV equations, it is possible to establish that the mass of the envelop contributes to less than 1% to the total mass of the star. Hence, Eq.(17) evaluated at $\rho_{\rm nuc}$ is a good approximation to the maximum mass that apply to any high density EOS:

$$M_{\rm max} \lesssim \sqrt{\frac{16}{243\pi\,\rho_{\rm nuc}}} \sim 5\,M_\odot \tag{18}$$

$\Rightarrow$ No (spherical) NS can exist above this mass.

The argument can be improved by adding the hypothesis that unknown EOS is causal, i.e. the speed of sound is less than the speed of light, $c_s^2 = dP/d\rho < 1$. This way, Roades&Ruffini (1974) found the improved, yet rather general, bound

$$M_{\rm max} \lesssim 3.2\,\sqrt{\frac{\rho_0}{\rho_{\rm nuc}}}\,M_\odot \tag{19}$$

Exercises

- Prove Buchdal's inequality (Follow Wald's book)
- Derive Eq.(14)
- Reproduce the calculations in Sec.(9.5) of Shapiro&Teukolsky book

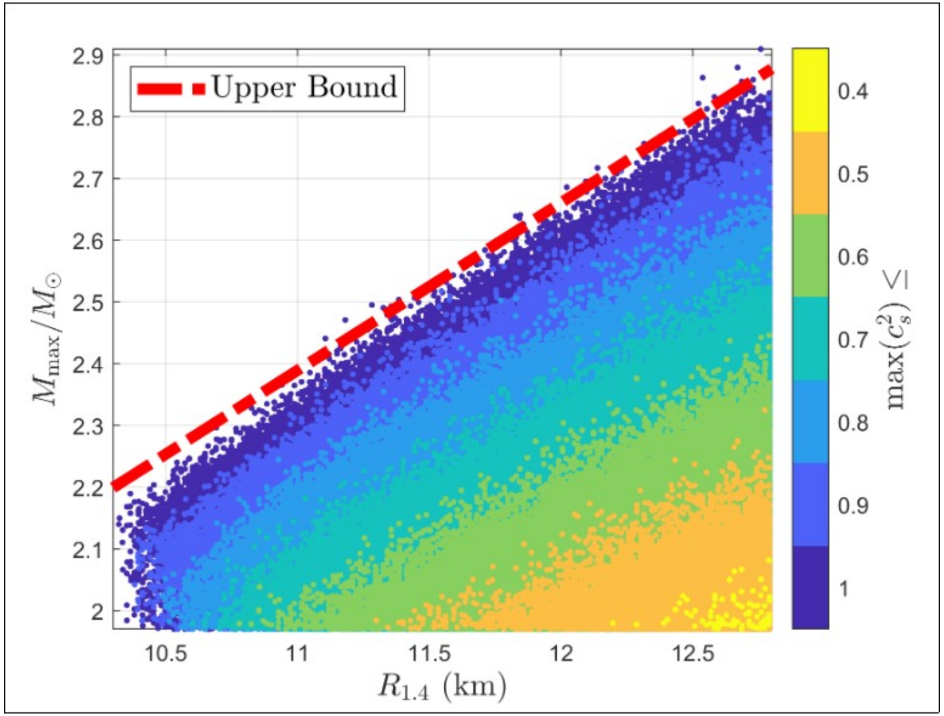
- GR & Causality
- No EOS assumptions
- ~2M phenomenological EOS

In general, we find that the upper bound on $M_{\rm max}$ can be very well approximated as

$$M_{\rm max} \leq \alpha(M) + \beta(M) R_M, \tag{2}$$

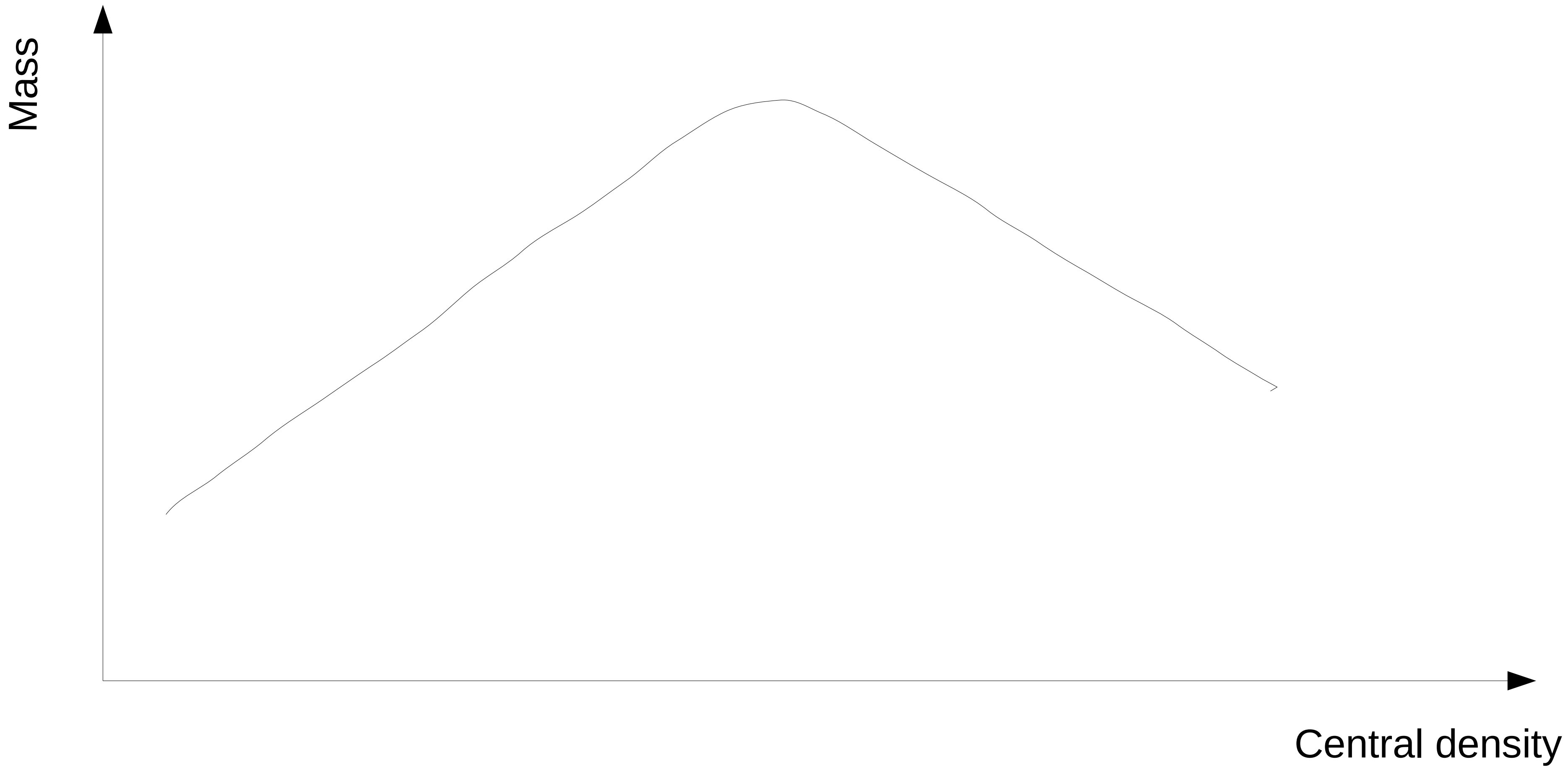
where R_M is the radius in km of a NS of gravitational mass M and

$$\begin{aligned} \alpha &= 0.45\,{\rm M}_\odot - 1.22\,M, \\ \beta &= -0.051\,{\rm M}_\odot\,{\rm km}^{-1} + 0.34\,M\,{\rm km}^{-1}. \end{aligned} \tag{3}$$



[Godzieba+ 2020]

Necessary condition for stability



NS pulsations

THE DYNAMICAL INSTABILITY OF GASEOUS MASSES APPROACHING THE SCHWARZSCHILD LIMIT IN GENERAL RELATIVITY

S. CHANDRASEKHAR
University of Chicago
Received May 11, 1964

ABSTRACT

In this paper the theory of the infinitesimal, baryon-number conserving, adiabatic, radial oscillations of a gas sphere is developed in the framework of general relativity. A variational base for determining the characteristic frequencies of oscillation is established. It provides a convenient method for obtaining sufficient conditions for the occurrence of dynamical instability. The principal result of the analysis is the demonstration that the Newtonian lower limit $\frac{4}{3}$, for the ratio of the specific heats γ, for insuring dynamical stability is increased by effects arising from general relativity; indeed, is increased to an extent that, so long as γ is finite, dynamical instability will intervene before a mass contracts to the limiting radius ($\geq 2.25\, GM/c^2$) compatible with hydrostatic equilibrium. Moreover, if γ should exceed $\frac{4}{3}$ only by a small amount, then dynamical instability will occur if the mass should contract to the radius

$$R_c = \frac{K}{\gamma - \frac{4}{3}} \frac{2GM}{c^2} \qquad (\gamma \to \tfrac{4}{3}),$$

where K is a constant depending, principally, on the density distribution in the configuration. The value of the constant K is explicitly evaluated for the homogeneous sphere of constant energy density and the polytropes of indices $n = 1, 2,$ and 3.

THE ASTROPHYSICAL JOURNAL, Vol. 149, September 1967

NON-RADIAL PULSATION OF GENERAL-RELATIVISTIC STELLAR MODELS. I. ANALYTIC ANALYSIS FOR $l \geq 2$*

KIP S. THORNE†
California Institute of Technology, Pasadena
AND
ALFONSO CAMPOLATTARO
University of California, Irvine
Received February 24, 1967

ABSTRACT

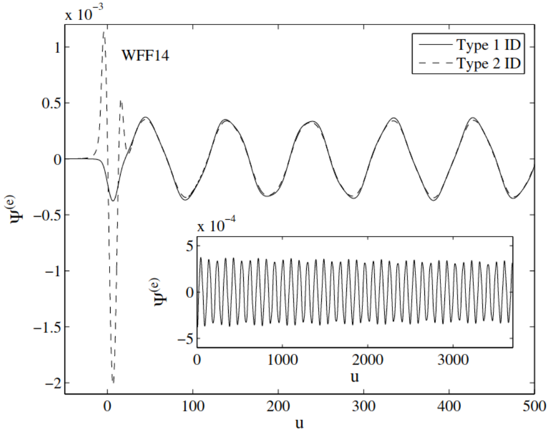
The theory of small, adiabatic, non-radial perturbations of a star away from hydrostatic equilibrium is developed within the framework of general relativity. The unperturbed equilibrium configuration is an arbitrary, non-rotating, general-relativistic stellar model. The departures from equilibrium are analyzed into tensorial spherical harmonics and then into complex normal modes with various mixtures of incoming and outgoing gravitational waves. A discussion is given of the expansion of real, physical pulsations with purely outgoing gravitational radiation in terms of the complex normal modes. Criteria are developed for stability against non-radial pulsations; and methods are devised for computing numerically the pulsation frequencies, eigenfunctions, and gravitational-radiation damping times of the stable, real quasi-normal modes of pulsation.

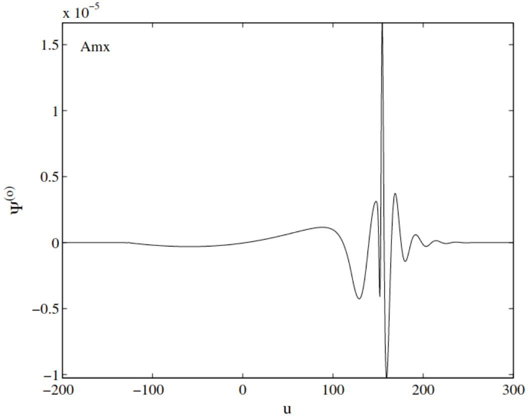
Fluid mode displacement:

$$\xi_n(r) \sim a_n(r) e^{i\omega_n t}$$

Stability: $\omega_n^2 < 0$

$$\frac{dM}{d\rho} = 0 \rightarrow \omega_n^2 = 0$$





Review: https://link.springer.com/article/10.12942/lrr-1999-2

Mon. Not. R. Astron. Soc. **299,** 1059–1068 (1998)

Towards gravitational wave asteroseismology

Nils Andersson[1,2] and Kostas D. Kokkotas[3,4]
[1] *Department of Physics, Washington University, St Louis, MO 63130, USA*
[2] *Institut für Theoretische Astrophysik, Universität Tübingen, D-72076 Tübingen, Germany*
[3] *Department of Physics, Aristotle University of Thessaloniki, Thessaloniki 54006, Greece*
[4] *Max-Planck-Institut für Gravitationsphysik, Schlaatzweg 1, 14473 Potsdam, Germany*

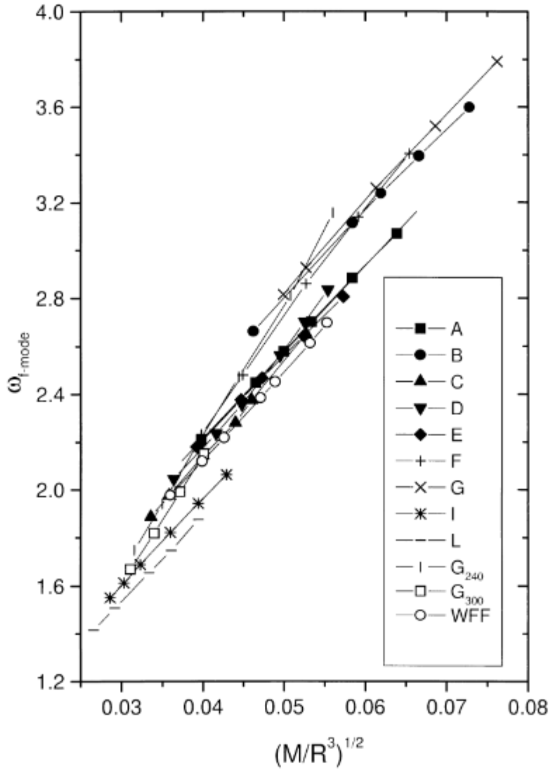


Figure 1. The numerically obtained *f* mode frequencies plotted as functions of the mean stellar density (M and R are in km and $\omega_{f\ \mathrm{mode}}$ in kHz).

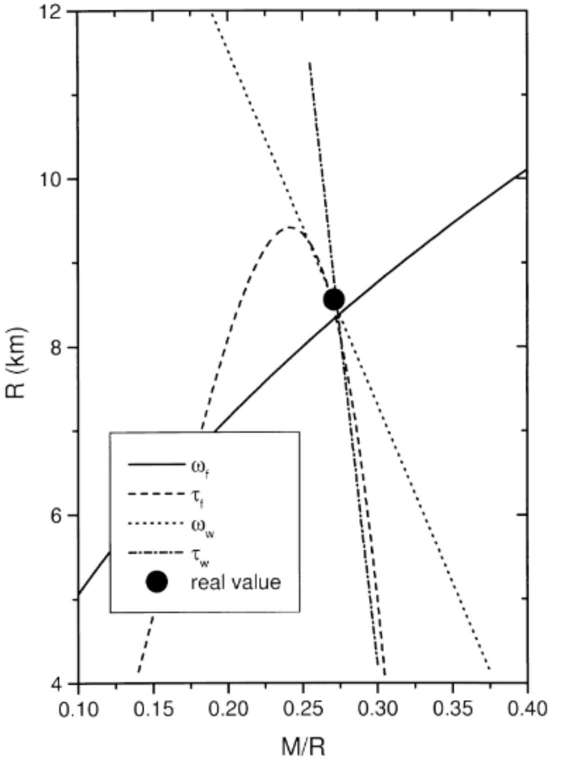
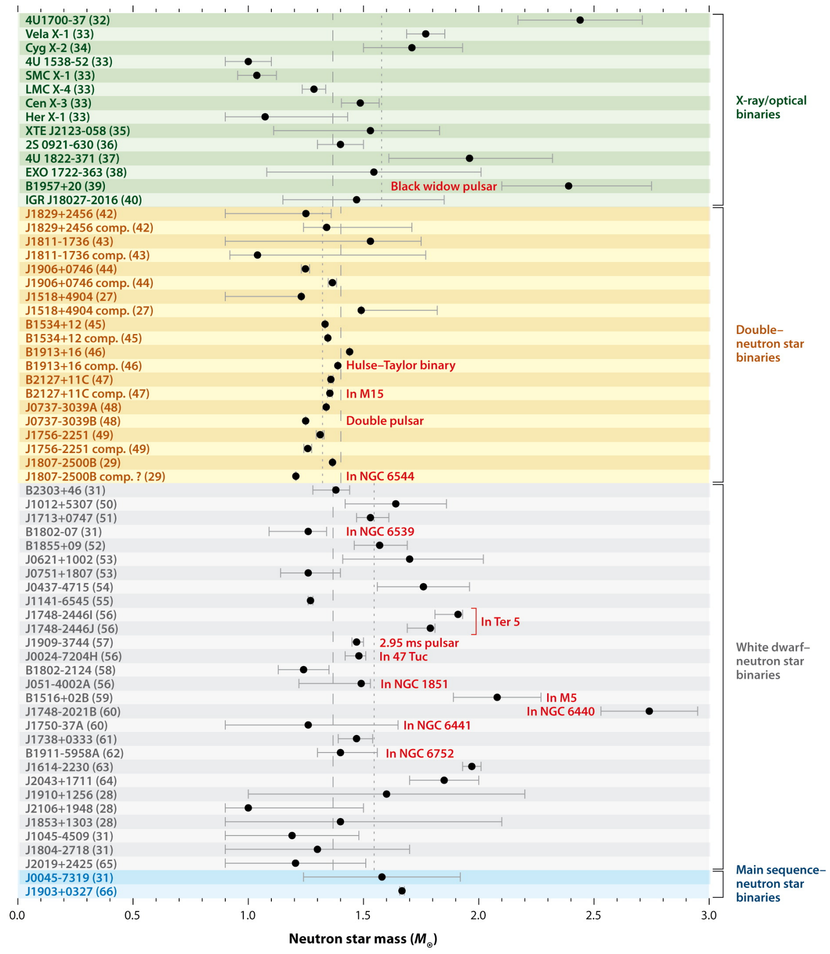


Figure 7. An illustration of how accurately the radius and the mass of a star can be inferred from detected mode data and our empirical relations.

Observed NS masses



Lattimer JM. 2012.
Annu. Rev. Nucl. Part. Sci. 62:485–515

A two-solar-mass neutron star measured using Shapiro delay

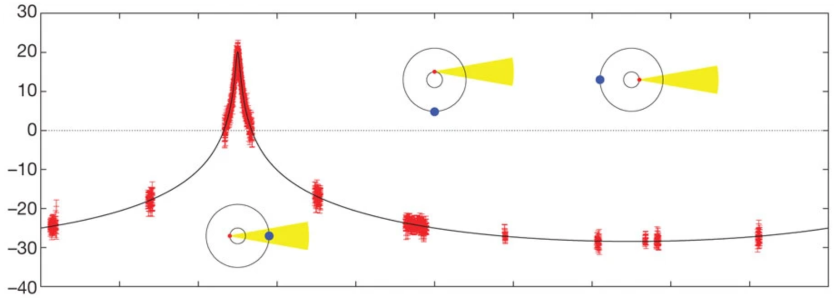
P. B. Demorest[1], T. Pennucci[2], S. M. Ransom[1], M. S. E. Roberts[3] & J. W. T. Hessels[4,5]

Neutron stars are composed of the densest form of matter known to exist in our Universe, the composition and properties of which are still theoretically uncertain. Measurements of the masses or radii of these objects can strongly constrain the neutron star matter equation of state and rule out theoretical models of their composition[1,2]. The observed range of neutron star masses, however, has hitherto been too narrow to rule out many predictions of 'exotic' non-nucleonic components[3–6]. The Shapiro delay is a general-relativistic increase in light travel time through the curved space-time near a massive body[7]. For highly inclined (nearly edge-on) binary millisecond radio pulsar systems, this effect allows us to infer the masses of both the neutron star and its binary companion to high precision[8,9]. Here we present radio timing observations of the binary millisecond pulsar J1614-2230[10,11] that show a strong Shapiro delay signature. We calculate the pulsar mass to be $(1.97 \pm 0.04)M_\odot$, which rules out almost all currently proposed[2–5] hyperon or boson condensate equations of state ($M_\odot$, solar mass). Quark matter can support a star this massive only if the quarks are strongly interacting and are therefore not 'free' quarks[12].

A Massive Pulsar in a Compact Relativistic Binary

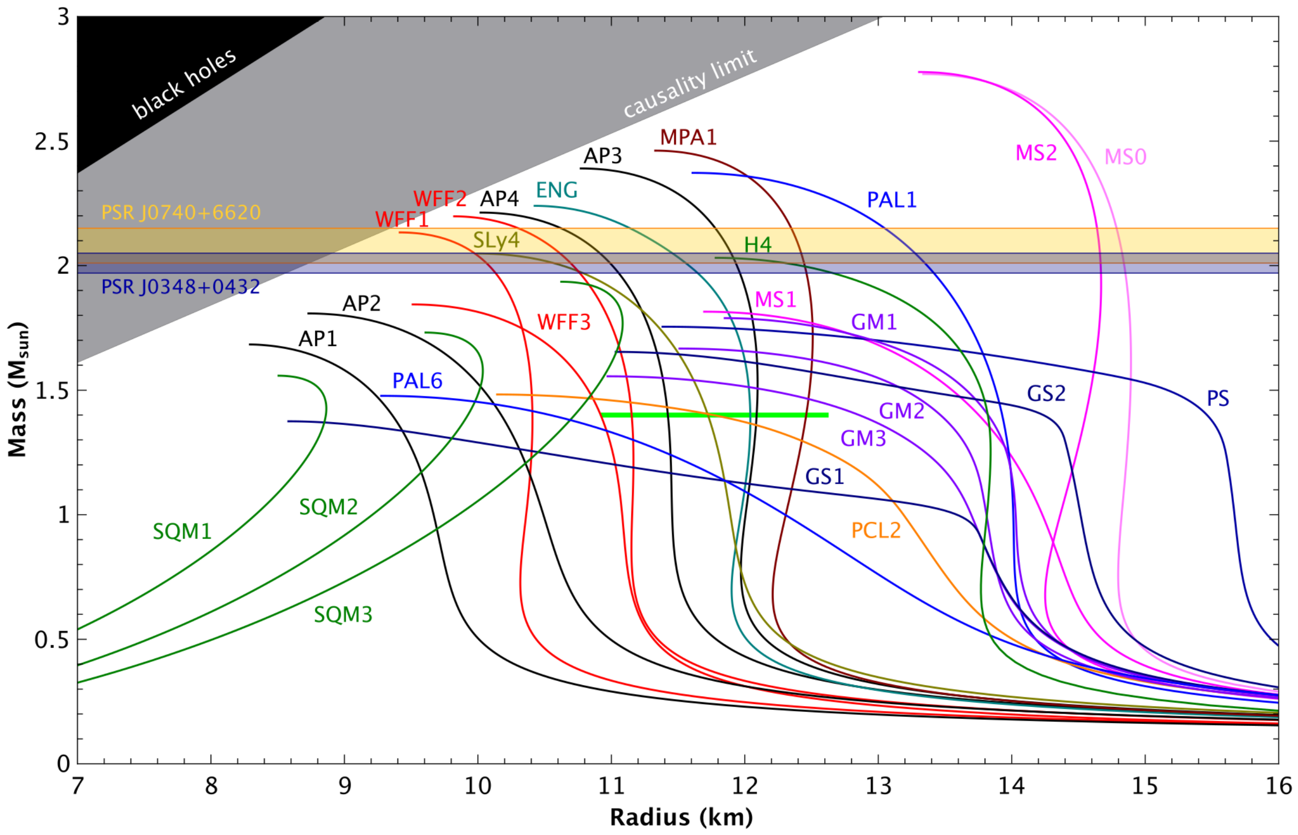
John Antoniadis,[1]* Paulo C. C. Freire,[1] Norbert Wex,[1] Thomas M. Tauris,[2,1] Ryan S. Lynch,[3] Marten H. van Kerkwijk,[4] Michael Kramer,[1,5] Cees Bassa,[5] Vik S. Dhillon,[6] Thomas Driebe,[7] Jason W. T. Hessels,[8,9] Victoria M. Kaspi,[3] Vladislav I. Kondratiev,[8,10] Norbert Langer,[2] Thomas R. Marsh,[11] Maura A. McLaughlin,[12] Timothy T. Pennucci,[13] Scott M. Ransom,[14] Ingrid H. Stairs,[15] Joeri van Leeuwen,[8,9] Joris P. W. Verbiest,[1] David G. Whelan[13]

Many physically motivated extensions to general relativity (GR) predict substantial deviations in the properties of spacetime surrounding massive neutron stars. We report the measurement of a 2.01 ± 0.04 solar mass ($M_\odot$) pulsar in a 2.46-hour orbit with a $0.172 \pm 0.003\ M_\odot$ white dwarf. The high pulsar mass and the compact orbit make this system a sensitive laboratory of a previously untested strong-field gravity regime. Thus far, the observed orbital decay agrees with GR, supporting its validity even for the extreme conditions present in the system. The resulting constraints on deviations support the use of GR-based templates for ground-based gravitational wave detectors. Additionally, the system strengthens recent constraints on the properties of dense matter and provides insight to binary stellar astrophysics and pulsar recycling.



Shapiro time delay: Timing residual as a function of pulsar's orbital phase

Mass-radius



From: https://www3.mpifr-bonn.mpg.de/staff/pfreire/NS_masses.html

GW170817: Measurements of Neutron Star Radii and Equation of State

The LIGO Scientific Collaboration and The Virgo Collaboration

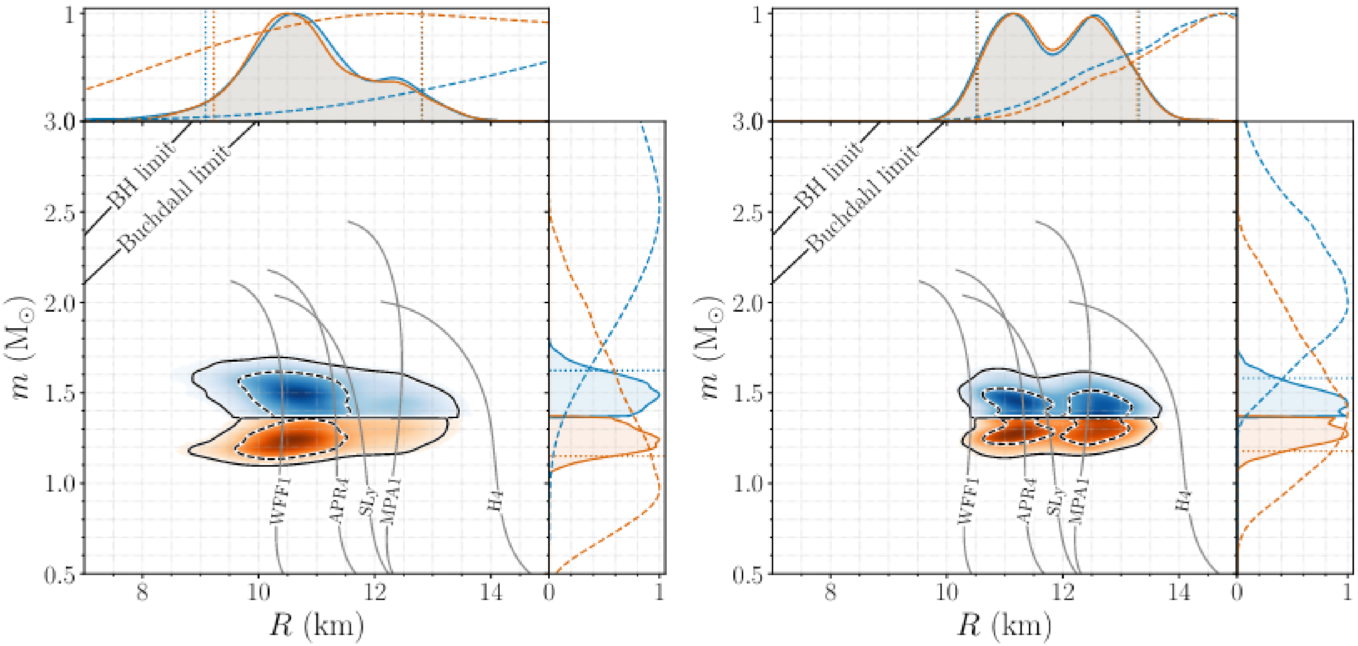
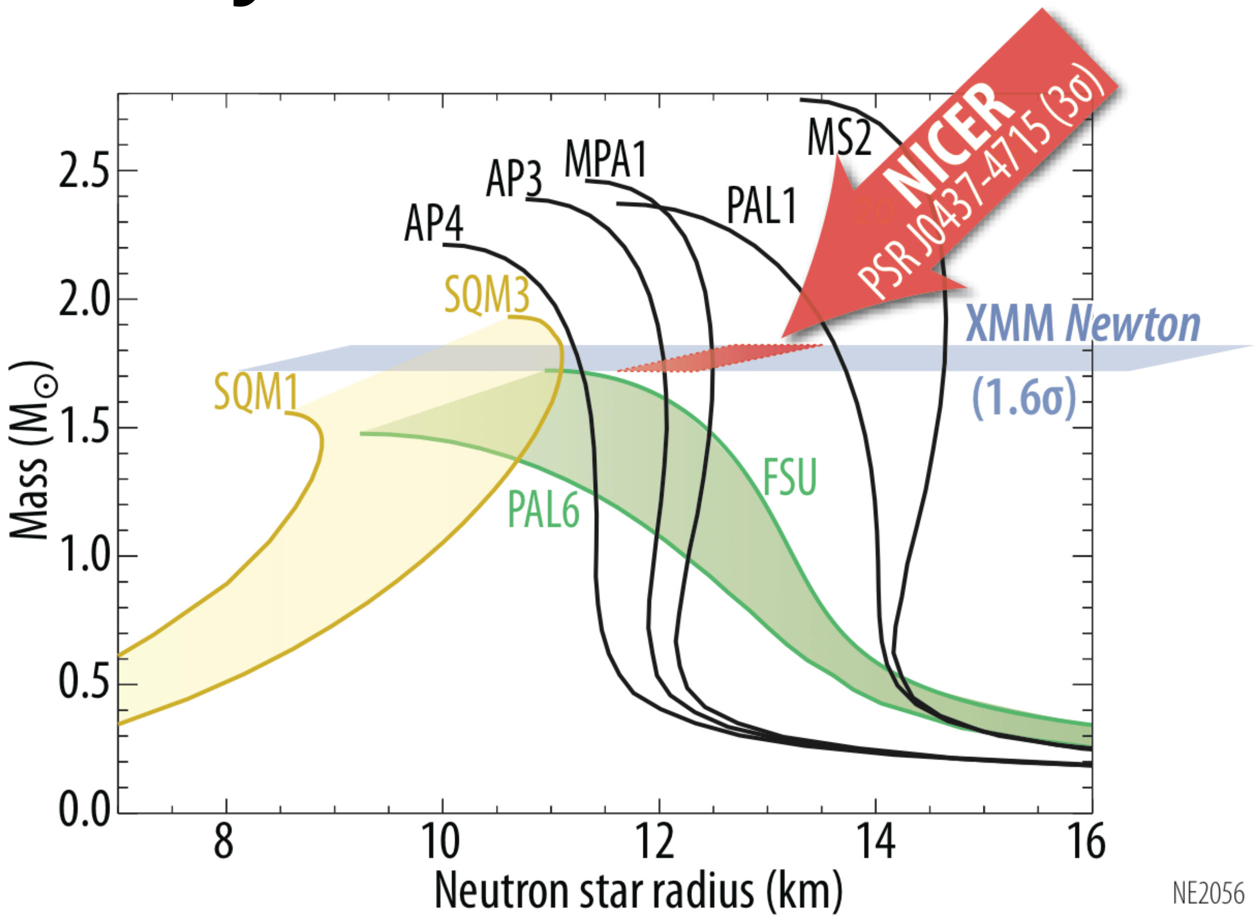


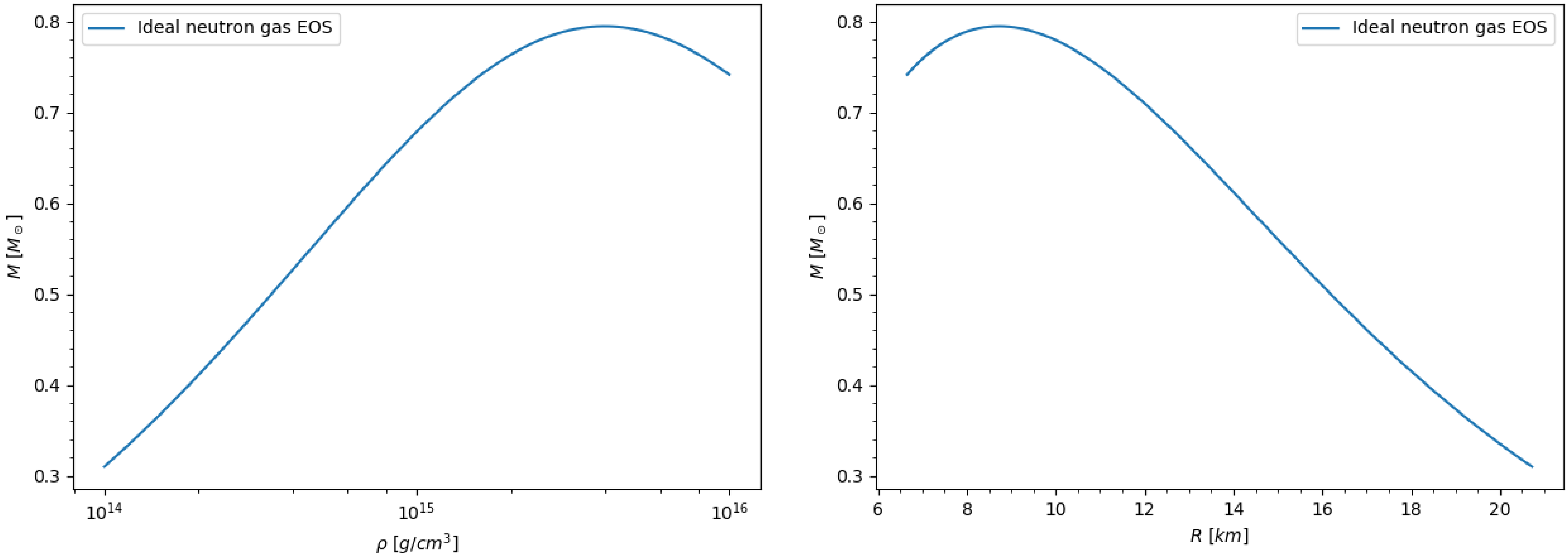
FIG. 3. Marginalized posterior for the mass m and areal radius R of each binary component using EOS-insensitive relations (left panel) and a parametrized EOS where we impose a lower limit on the maximum mass of 1.97 $M_\odot$ (right panel). The top blue (bottom orange) posterior corresponds to the heavier (lighter) NS. Example mass-radius curves for selected EOSs are overplotted in gray. The lines in the top left denote the Schwarzschild BH ($R = 2m$) and Buchdahl ($R = 9m/4$) limits. In the one-dimensional plots, solid lines are used for the posteriors, while dashed lines are used for the corresponding parameter priors. Dotted vertical lines are used for the bounds of the 90% credible intervals.

X-ray measurements



NICER https://heasarc.gsfc.nasa.gov/docs/nicer/

Neutron-gas stars



NS EOS

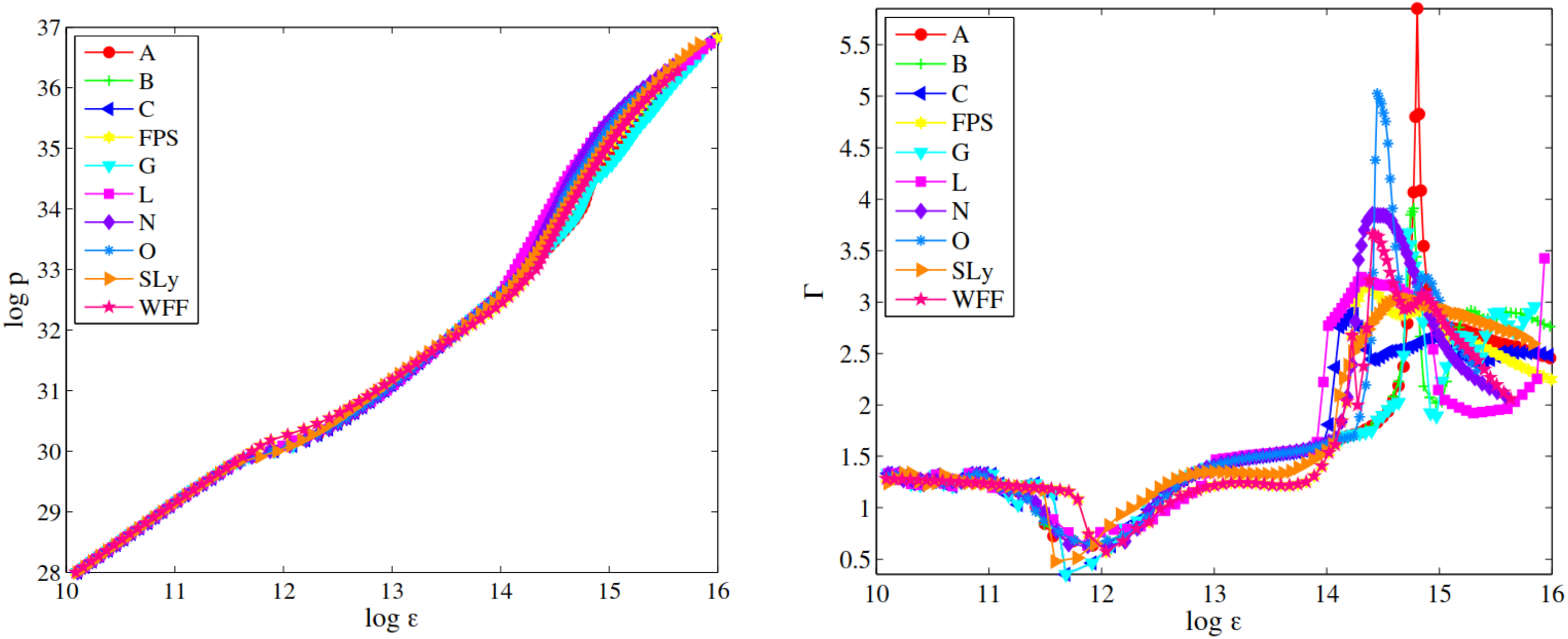
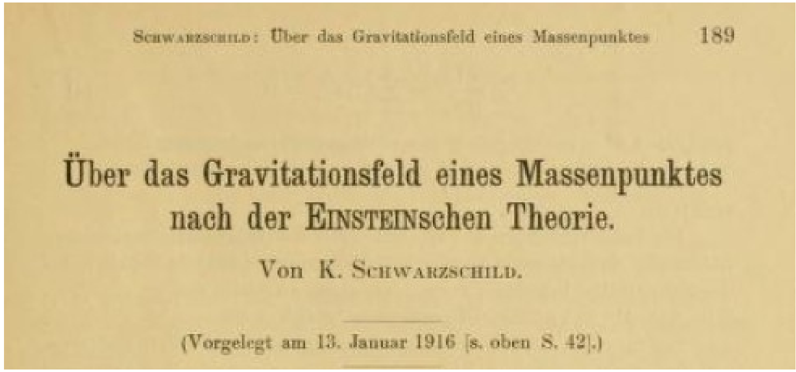


FIG. 1: Pressure (top) and the adiabatic exponent (bottom) as a function of the total energy density for various EOS. Notice here we are using cgs units.

Black holes

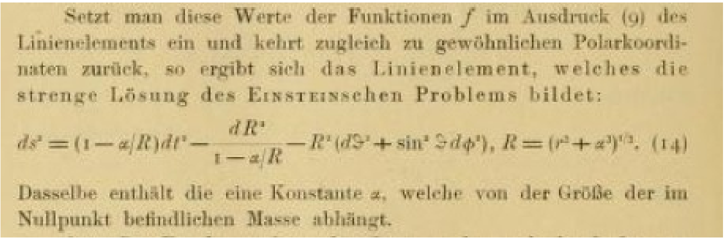


SCHWARZSCHILD: Über das Gravitationsfeld eines Massenpunktes 189

Über das Gravitationsfeld eines Massenpunktes nach der EINSTEINschen Theorie.

Von K. SCHWARZSCHILD.

(Vorgelegt am 13. Januar 1916 [s. oben S. 42].)



Setzt man diese Werte der Funktionen f im Ausdruck (9) des Linienelements ein und kehrt zugleich zu gewöhnlichen Polarkoordinaten zurück, so ergibt sich das Linienelement, welches die strenge Lösung des EINSTEINschen Problems bildet:

$$ds^2 = (1-\alpha/R)dt^2 - \frac{dR^2}{1-\alpha/R} - R^2(d\vartheta^2 + \sin^2\vartheta d\phi^2),\ R = (r^3+\alpha^3)^{1/3}. \quad (14)$$

Dasselbe enthält die eine Konstante α, welche von der Größe der im Nullpunkt befindlichen Masse abhängt.

ASTRONOMY   SCIENCE NEWS LETTER for January 18, 1964

"Black Holes" in Space

The heavy densely packed dying stars that speckle space may help determine how matter behaves when enclosed in its own gravitational field—By Ann Ewing

➤ SPACE may be peppered with "black holes."

This was suggested at the American Association for the Advancement of Science meeting in Cleveland by astronomers and physicists who are experts on what are called degenerate stars.

Degenerate stars are not Hollywood types with low morals. They are dying stars, or white dwarfs, and make up about 10% of all stars in the sky.

The faint light they emit comes from the little heat left in their last stages of life. It is not known how a star quietly declines to become a white dwarf.

Degenerate stars are made of densely packed electrons and nuclei, or cores of atoms. They are so dense that a thimbleful of their matter weighs a ton.

Some such stars are predicted in theory to have a density of one million tons per thimbleful. When this happens, the star is essentially made of neutrons and strange particles.

Because a degenerate star is so dense, its gravitational field is very strong. According to Einstein's general theory of relativity, as mass is added to a degenerate star a sudden collapse will take place and the intense gravitational field of the star will close in on itself.

Such a star then forms a "black hole" in the universe.

Modern tools, such as telescopes on an orbiting space platform, may be used to detect such black holes and to help determine how matter behaves when it is enclosed by its own gravitational field.

The light from the most famous white dwarf star, Sirius B, a companion to Sirius—which is the brightest star in the heavens visible from earth—has been captured using the 200-inch telescope atop Mt. Palomar. This was done as part of a program to study at least 20 white dwarfs.

Preliminary analysis of the light from Sirius B indicates that it has an effective temperature of 16,800 degrees Kelvin, or 30,000 degrees Fahrenheit. Its radius can be calculated from the temperature, and is only nine-thousandths that of the sun.

The star must therefore consist mainly of helium or heavier elements.

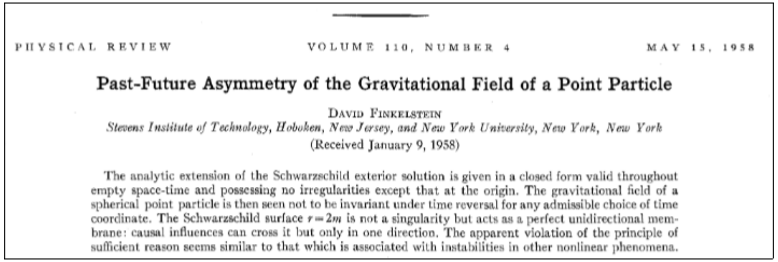
The speakers at the symposium were Drs. A. G. W. Cameron of the National Aeronautics and Space Administration's Goddard Institute for Space Studies, New York; Charles Misner of the University of Maryland; Volker Weidemann, Physikalisch-Technische Bundesanstalt, Braunschweig, Germany, and J. B. Oke of California Institute of Technology. The symposium was arranged by Dr. Hong-yee Chiu of the Goddard Institute for Space Studies.

• Science News Letter, 85:39 Jan. 18, 1964

SPACE

Original paper: https://archive.org/details/sitzungsberichte1916deutsch/page/188/mode/2up?view=theater

Translation: https://arxiv.org/abs/physics/9905030



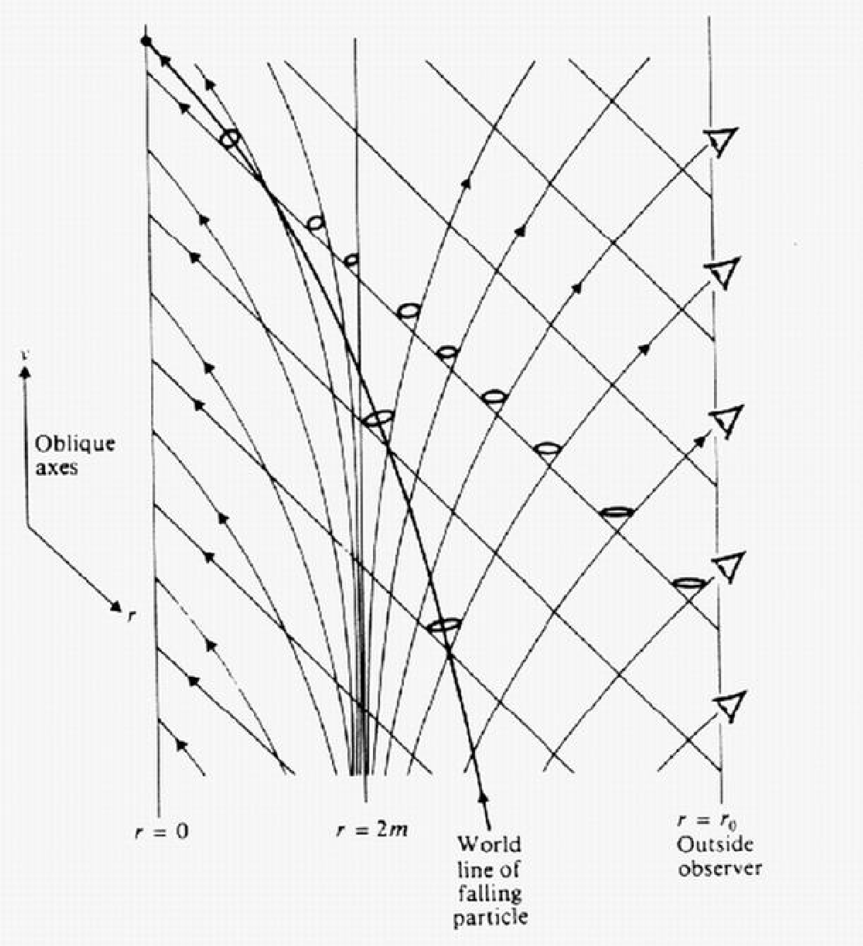
PHYSICAL REVIEW   VOLUME 110, NUMBER 4   MAY 15, 1958

Past-Future Asymmetry of the Gravitational Field of a Point Particle

DAVID FINKELSTEIN
Stevens Institute of Technology, Hoboken, New Jersey, and New York University, New York, New York
(Received January 9, 1958)

The analytic extension of the Schwarzschild exterior solution is given in a closed form valid throughout empty space-time and possessing no irregularities except that at the origin. The gravitational field of a spherical point particle is then seen not to be invariant under time reversal for any admissible choice of time coordinate. The Schwarzschild surface $r=2m$ is not a singularity but acts as a perfect unidirectional membrane: causal influences can cross it but only in one direction. The apparent violation of the principle of sufficient reason seems similar to that which is associated with instabilities in other nonlinear phenomena.

- Eddington (1924)
- Lemaitre (1933)
- Finkelstein (1958) "*a perfect unidirectional membrane: causal influences can cross it in only one direction*"



Clock's Redshift:

$$\frac{ds_\infty^2}{ds_r^2} = \frac{\sqrt{-g_{00}(\infty)}}{\sqrt{-g_{00}(r)}} = \frac{1}{\sqrt{1-\frac{2GM}{c^2 r}}}$$

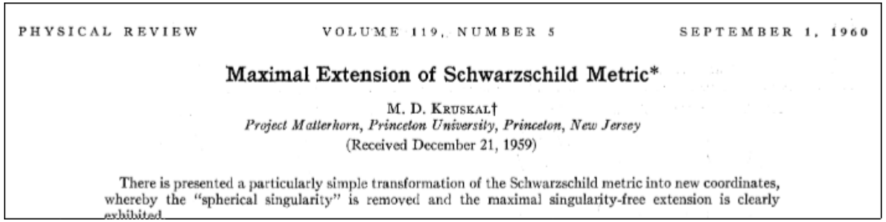
One way membrane:

Light cones tilt for $r < 2M$,

future directed paths are in the direction

of $r = 0$ (true singularity).

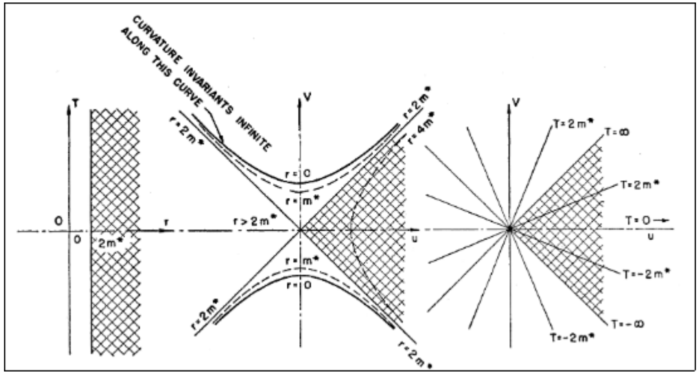
$r = 2M$ is a null surface called **event horizon**

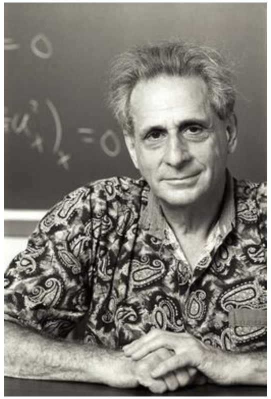

PHYSICAL REVIEW VOLUME 119, NUMBER 5 SEPTEMBER 1, 1960

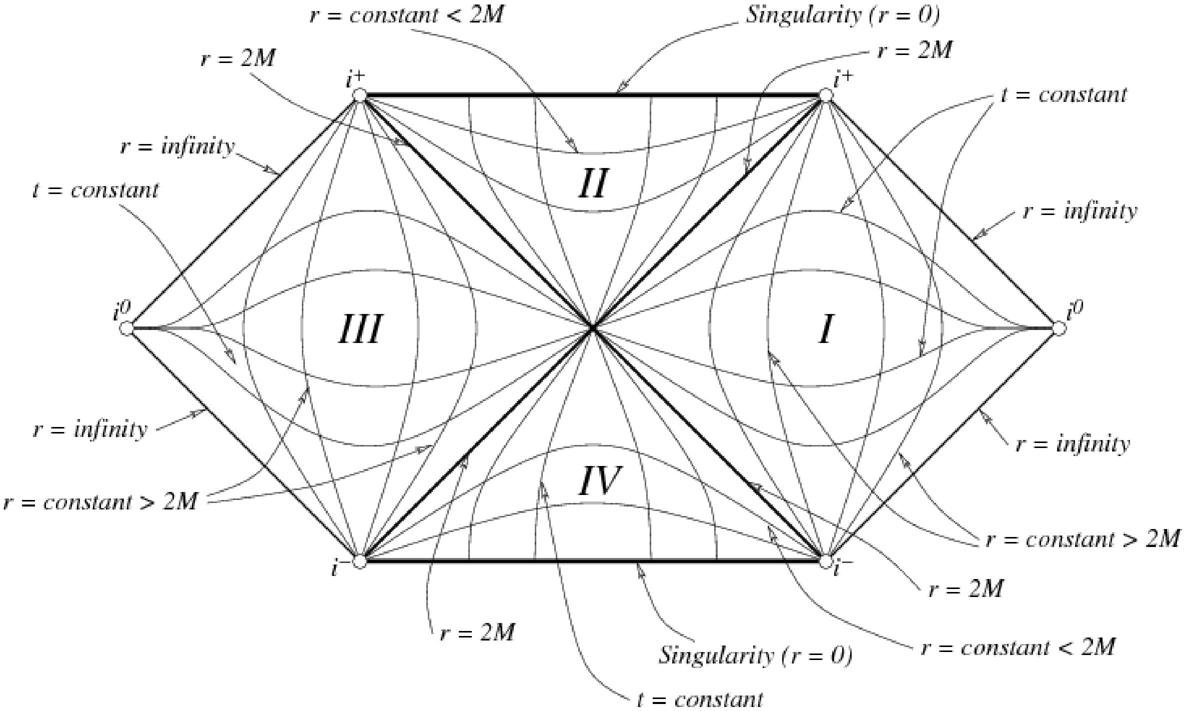
Maximal Extension of Schwarzschild Metric*

M. D. Kruskal†
Project Matterhorn, Princeton University, Princeton, New Jersey
(Received December 21, 1959)

There is presented a particularly simple transformation of the Schwarzschild metric into new coordinates, whereby the "spherical singularity" is removed and the maximal singularity-free extension is clearly exhibited.

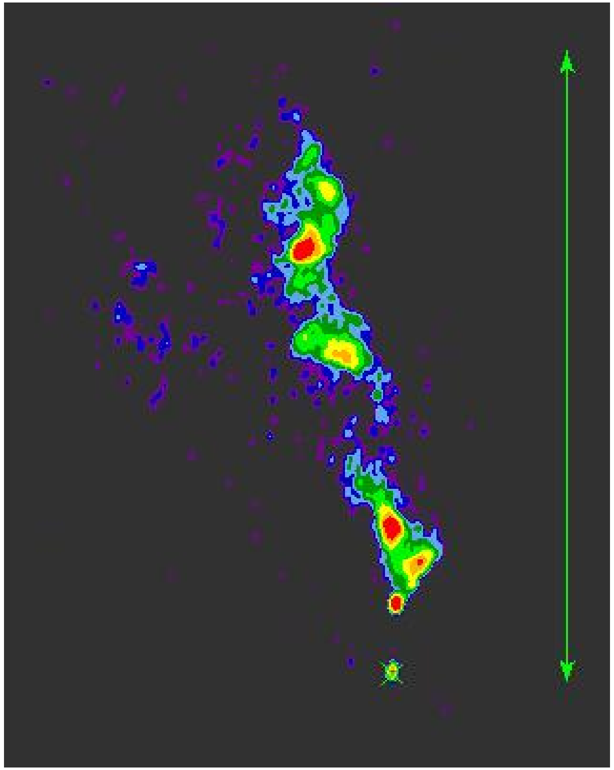




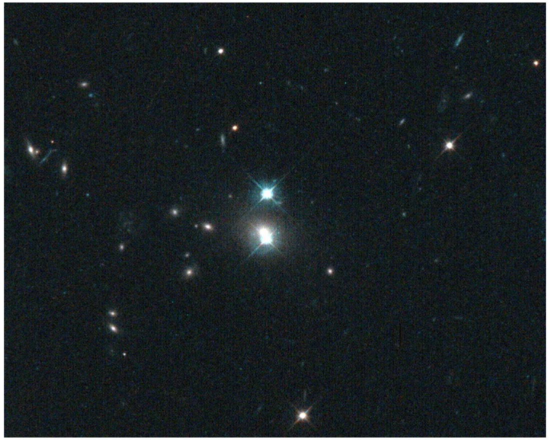


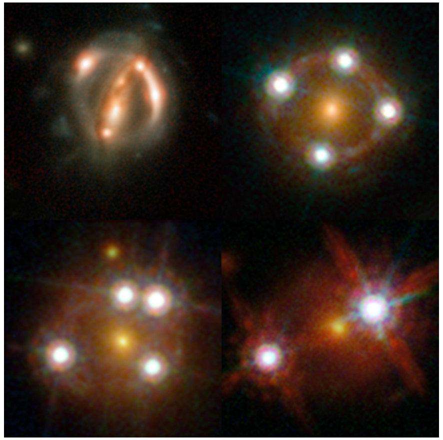
Quasar

C348 @ 1.4GHz



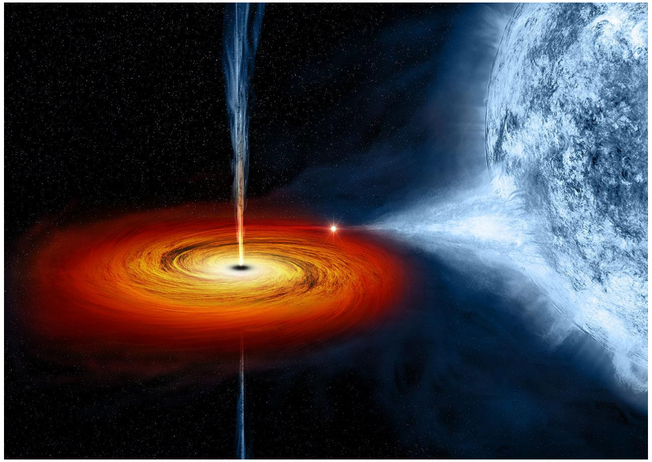
- 50s Radio sources of small size
- 60s Optical counterparts w\ High redshift (z ~ 7)
- Very luminous & extra-galactic? (> nuclear fusion, supernovae)
- 1964 Salpeter&Zeldovich: **Supermassive BH + accretion disk**
- Confirmed by
  - X ray observations of BH (*next slide*)
  - 1971 Peterson and Gunn: Galaxies containing quasars showed the same redshift as the quasars
  - 1979 Walsh,Carswell&Weyman: Grav. Lensing

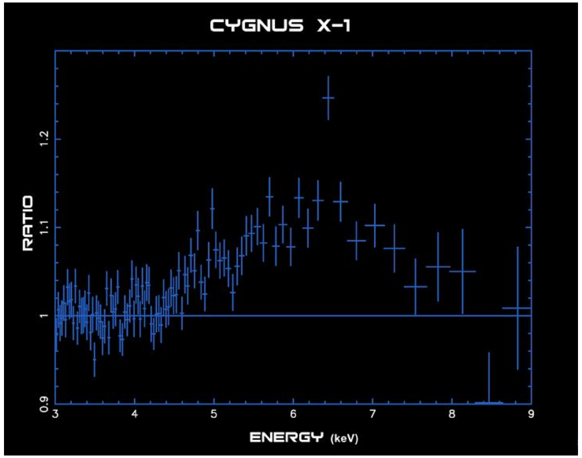




X-ray astronomy

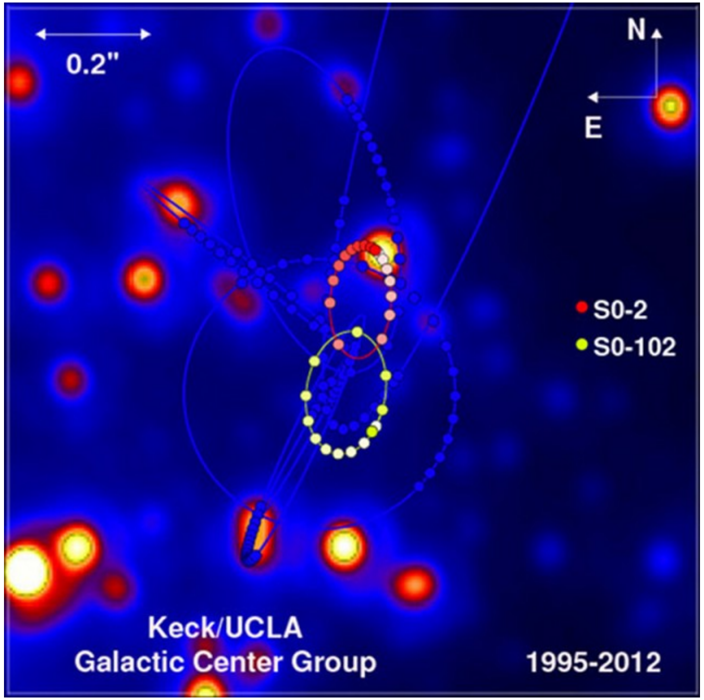
- Hot gases at T ~ 1,000,000K emit X-ray
- 1962 Scorpius X-1
  - Strongest X-ray source together the Sun.
  - Low-Mass-X-ray binary
  - 1.4M_{Sun} NS + 0.42 star
- 1964 Cygnus X-1
  - High-Mass-X-ray binary
  - 14.8M_{Sun} BH + 20-40M_{Sun} supergiant star

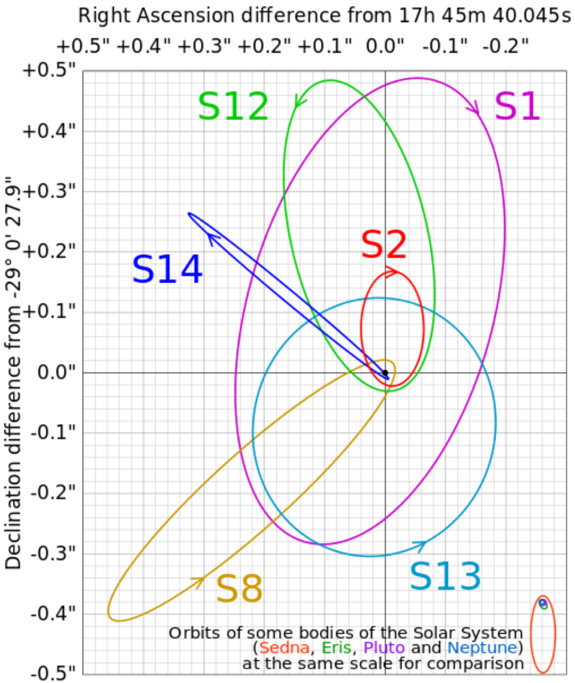




R.Giacconi (Nobel Prize 2002)

Sagittarius* A



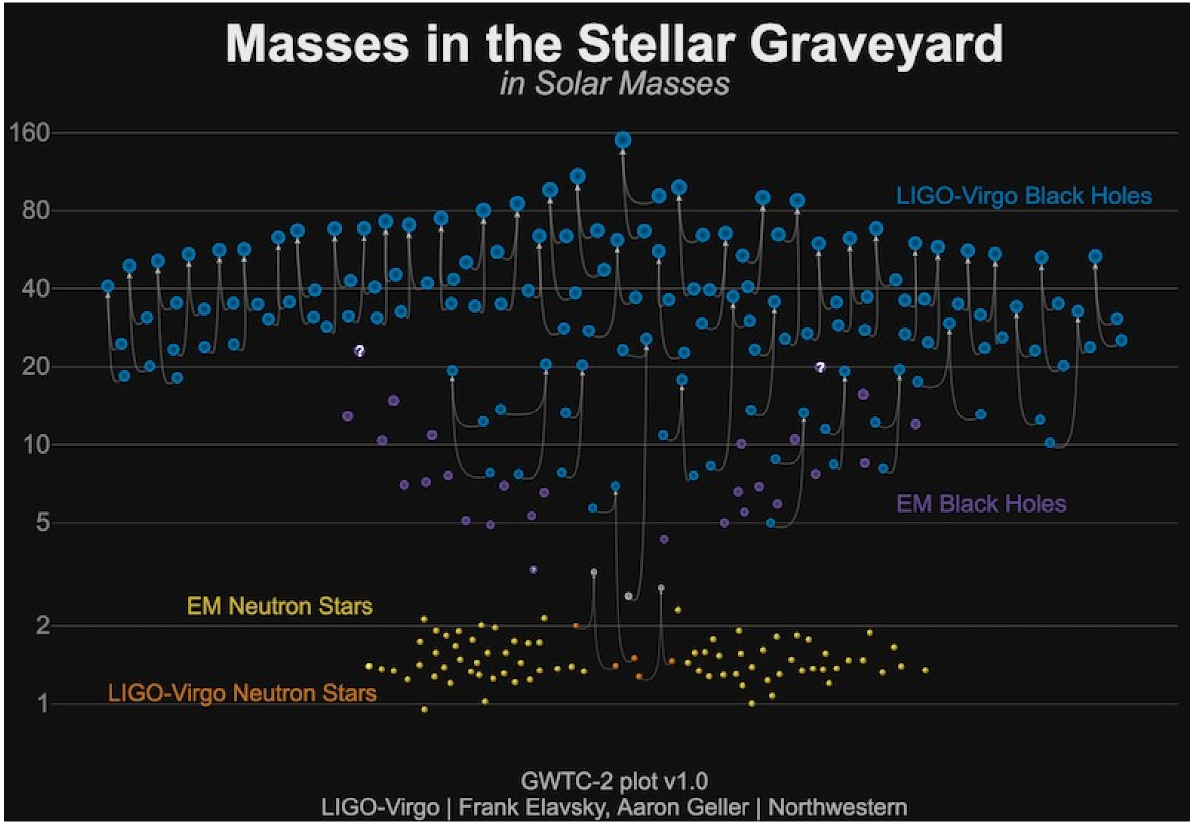


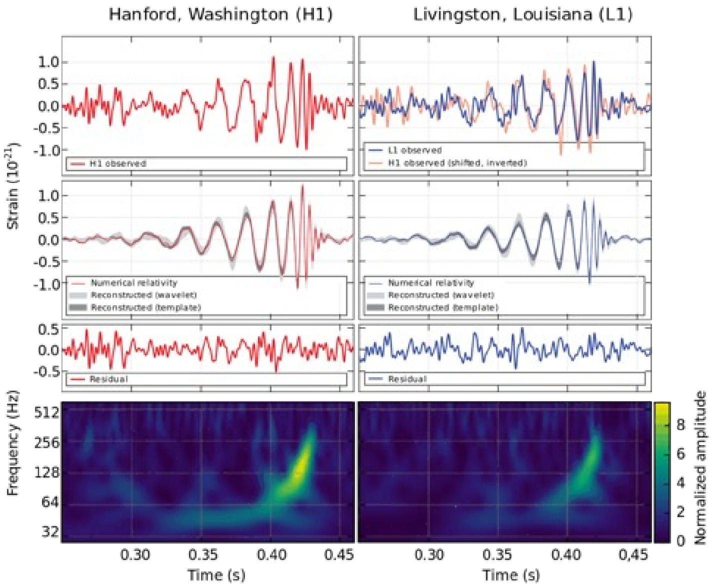
Galaxy center; Orbits' speed ~ 2% c

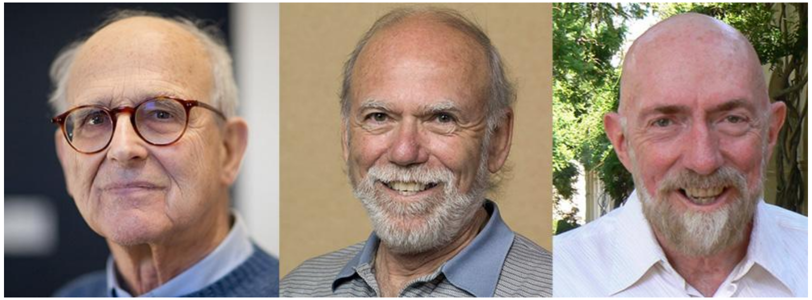
Mass ~ 4 million M_{Sun} ! => Supermassive BH

Gravitational-wave observations

Since 2015, LIGO-Virgo observations







Weiss,Barish,Thorne Nobel Prize 2017

Gravitational collapse

SEPTEMBER 1, 1939 PHYSICAL REVIEW VOLUME 56

On Continued Gravitational Contraction

J. R. Oppenheimer and H. Snyder
University of California, Berkeley, California
(Received July 10, 1939)

When all thermonuclear sources of energy are exhausted a sufficiently heavy star will collapse. Unless fission due to rotation, the radiation of mass, or the blowing off of mass by radiation, reduce the star's mass to the order of that of the sun, this contraction will continue indefinitely. In the present paper we study the solutions of the gravitational field equations which describe this process. In I, general and qualitative arguments are given on the behavior of the metrical tensor as the contraction progresses: the radius of the star approaches asymptotically its gravitational radius; light from the surface of the star is progressively reddened, and can escape over a progressively narrower range of angles. In II, an analytic solution of the field equations confirming these general arguments is obtained for the case that the pressure within the star can be neglected. The total time of collapse for an observer comoving with the stellar matter is finite, and for this idealized case and typical stellar masses, of the order of a day; an external observer sees the star asymptotically shrinking to its gravitational radius.

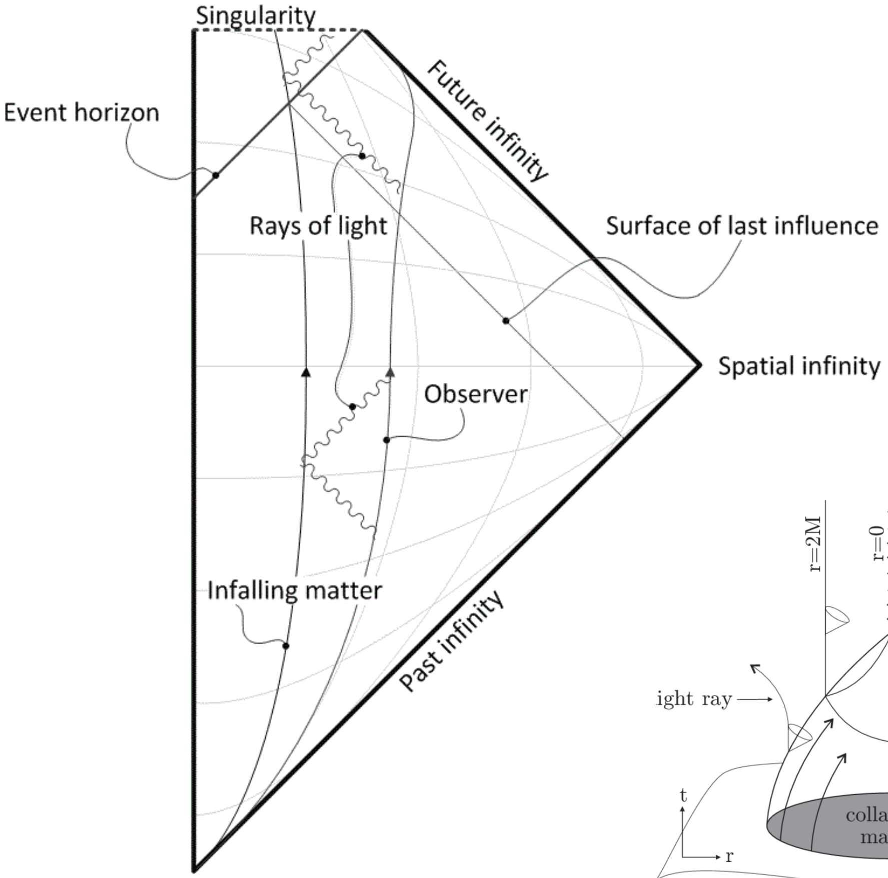
GRAVITATIONAL COLLAPSE AND SPACE-TIME SINGULARITIES

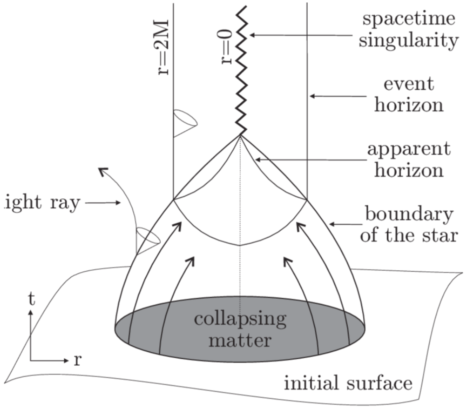
Roger Penrose
Department of Mathematics, Birkbeck College, London, England
(Received 18 December 1964)

The discovery of the quasistellar radio sources has stimulated renewed interest in the question of gravitational collapse. It has been suggested by some authors[1] that the enormous amounts of energy that these objects apparently emit may result from the collapse of a mass of the order of (10^6-10^8)$M_\odot$ to the neighborhood of its Schwarzschild radius, accompanied by a violent release of energy, possibly in the form of gravitational radiation. The detailed mathematical discussion of such situations is difficult since the full complexity of general relativity is required. Consequently, most exact calculations concerned with the implications of gravitational collapse have employed the simplifying assumption of spherical symmetry. Unfortunately, this precludes any detailed discussion of gravitational radiation—which requires at least a quadripole structure.

measured by local comoving observers, the body passes within its Schwarzschild radius $r=2m$. (The densities at which this happens need not be enormously high if the total mass is large enough.) To an outside observer the contraction to $r=2m$ appears to take an infinite time. Nevertheless, the existence of a singularity presents a serious problem for any complete discussion of the physics of the interior region.

The question has been raised as to whether this singularity is, in fact, simply a property of the high symmetry assumed. The matter collapses radially inwards to the single point at the center, so that a resulting space-time catastrophe there is perhaps not surprising. Could not the presence of perturbations which destroy the spherical symmetry alter the situation drastically? The recent rotating





Birkhoff's Theorem

Theorem 1. *Birkhoff (1923). The Schwarzschild metric is the unique vacuum solution in spherical symmetry.*

Sketch of the proof.

i. Any spherically symmetric spacetime (three spacelike rotational Killing vectors) can be foliated in 2-spheres

ii. The most general form of the metric is

$$g = -e^{2\phi(t,r)}dt^2 + e^{2\lambda(t,r)}d^2r + r^2 d^2\Omega \tag{1}$$

iii. Use EFE to "eliminate" the time dependence

Corollary 2. *Any spherically symmetric vacuum spacetime is static.*

Physically, the staticity result can be understood as the absence of gravitational monopole radiation (analogous to the fact that the Coulomb solution is the only spherically symmetric solution of Maxwell equations in vacuum).

For example the exterior spacetime of a gravitationally collapsing spherical body is static always given by the static Schwarzschild metric.

Orbits

Geodesics of photons and particles in the Schwazrschild metric can be analyzed introducing the constants of motions associated to each Killing vector K^α of the spacetime. Exactly as in the Newtonian problem the motion is on a plane and the relevant equations are

$$-T_\alpha \frac{dx^\alpha}{d\tau} = \underbrace{\left(1 - \frac{2M}{r}\right)}_{=:A(r)} \dot{t} = \text{const} =: E \quad \text{energy} \tag{2}$$

$$\Phi_\alpha \frac{dx^\alpha}{d\tau} = r^2 \dot{\phi} = p_\phi = \text{const} =: L \quad \text{angular momentum} \tag{3}$$

$$-s = g_{\alpha\beta} \frac{dx^\alpha}{d\tau} \frac{dx^\beta}{d\tau} \tag{4}$$

where $s = 0, 1$ for photons and unit-mass test particles respectively. The key equation resulting from the ones above is remarkably simple:

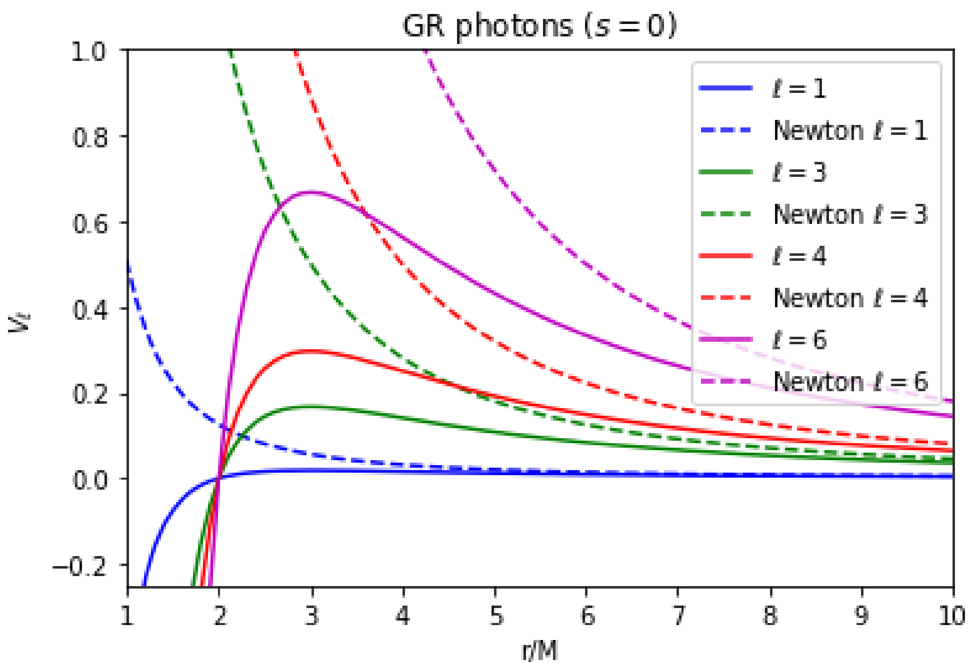
$$\dot{r} + V = E^2 \tag{5}$$

with the potential

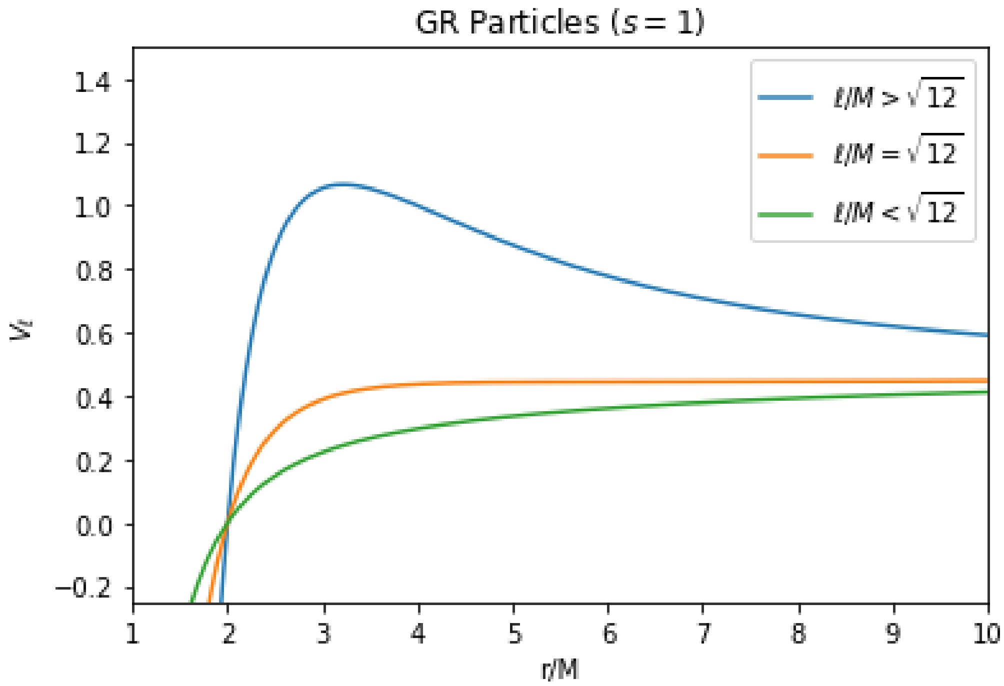
$$V_L := A(r)\left(s + \frac{L^2}{r^2}\right) = s - s\frac{2M}{r} + \frac{L^2}{r^2} - \frac{2ML^2}{r^3} \tag{6}$$

This result is analogous to the Newtonian motion in a central potential plus a GR term $\sim r^{-3}$.

Orbits



GR photons ($s = 0$)
$\ell = 1$
Newton $\ell = 1$
$\ell = 3$
Newton $\ell = 3$
$\ell = 4$
Newton $\ell = 4$
$\ell = 6$
Newton $\ell = 6$
V_ℓ
r/M



GR Particles ($s = 1$)
$\ell/M > \sqrt{12}$
$\ell/M = \sqrt{12}$
$\ell/M < \sqrt{12}$
V_ℓ
r/M

The analysis of the orbits is thus performed by analyzing the stationary points of the potential (Note $\dot{r}^2 = E^2 - V_L \geq 0$):

$$0 = \frac{dV_L}{dr} = sMr^2 - L^2 r + 3ML^2 \qquad \frac{d^2 V_L}{dr^2} = 2sMr - L^2 \tag{7}$$

Short summary of relevant feats:

Photons (s=0)

- The potential has maximum at $r = 3M$ (for $L > 0$), called *light ring* with energy $E_{\mathrm{LR}} = \sqrt{V(3M)} = \sqrt{L^2/(27M^2)}$
- The light ring corresponds to an *unstable circular orbit*
- Incoming photons with $E > E_{\mathrm{LR}}\,(E < E_{\mathrm{LR}})$ continue to $r = 2M$ and below (hit a turning point at a minimum radius and reverse the trajectory)

Particles (s=1)

- The potential has extrema at $r_\pm = L^2 \pm \sqrt{L^2(L^2 - 12M^2)}$ with energies $E_\pm$
- The values $r_\pm$ correspond to an unstable and a stable circular orbit respectively
- $r = r_+ = r_- = 6M$ is the last stable orbit (LSO) or innermost stable circular orbits (ISCO)
- Incoming particles with $E > E_-$ $(E < E_+)$ continue to $R = 2M$ (hit a turning point and reverse)
- Particles with $E_- < E < E_+$ move on bound orbits (not necessarily closed; precession)

The simplest relativistic two-body problem

Imagine a small but finite mass on a circular orbit around a nonrotating black hole. The emission of gravitational radiation determines a deviation from geodesic motion. If initially $r \gg 2M$, the emission timescale is much longer than the orbital period and one can approximate the dynamics as a sequence of circular orbits with progressively smaller radius and higher frequency (*adiabatic approximation*). While at some point the adiabatic approximation will break, we can still analyze the motion and make some predictons/estimates.

The orbital radius will continue decreasing to the LSO. Below that point, no stable circular orbit is possible and the particle will fall to $r = 2M$ and then down to $r = 0$.

$\Rightarrow$ The "two bodies" collide and merge!

The orbital frequency of the LSO is easily found from the angular momentum value at $r_+ = 6M$:

$$\Omega^2 = \frac{L^2}{r^4} = \frac{M}{r^2(r - 3M)} = 6^{-3/2} M^{-2} \tag{8}$$

or $(M\Omega)^2 = 6^{-3/2}$, the corresponding gravitational frequency is twice this value and provides an estimate of the merger frequency of the binary. Similarly, the energy of the LSO is

$$E = \frac{r - 2M}{[r(r - 3M)]^{1/2}} = \sqrt{\frac{8}{9}} \tag{9}$$

Thus, the energy emitted in gravitational waves (per unit mass) is $1 - E \approx 0.06$.

Exercises

1. Derive the formulas used above to discuss the orbits.
2. Derive the Hamiltonian of particles [Hint: Start from circular orbits]
3. Estimate the gravitational-wave merger frequency of a binary neutron star made of two euqal-masses neutron stars of $1.4M_\odot$. Comment about the result.
4. Estimate the gravitational-wave merger frequency of an equal-mass binary black hole system of stellar-mass black holes of $30M_\odot$ and supermassive black holes of $10^6 M_\odot$.
5. The correct result of the previous exercise is $2M\Omega \simeq 0.36$. Can you say why it holds for both cases (i.e. why there is a trivial mass scale)?

Perturbations & Stability

PHYSICAL REVIEW VOLUME 108, NUMBER 4 NOVEMBER 15, 1957

Stability of a Schwarzschild Singularity

TULLIO REGGE, *Istituto di Fisica della Università di Torino, Torino, Italy*

AND

JOHN A. WHEELER, *Palmer Physical Laboratory, Princeton University, Princeton, New Jersey*

(Received July 15, 1957)

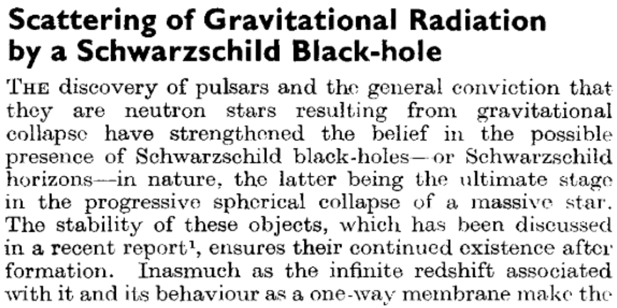
It is shown that a Schwarzschild singularity, spherically symmetrical and endowed with mass, will undergo small vibrations about the spherical form and will therefore remain stable if subjected to a small nonspherical perturbation.

$$-\Psi^{\ell m}_{tt} + \Psi^{\ell m}_{r_* r_*} - V_\ell(r_*(r))\Psi^{\ell m} = S_{\ell m}$$

$$r_* = r + R_s \ln\left(\frac{r}{R_s} - 1\right)$$

$$h_+ - i h_\times = \frac{G}{c^4 r}\sum_{\ell=2}\sum_{m=-\ell}^{\ell}\sqrt{\frac{(\ell+2)!}{(\ell-2)!}}\,{}^{(-2)}Y_{\ell m}(\theta,\varphi)\left(\Psi^{(e)}_{\ell m} + i\Psi^{(o)}_{\ell m}\right) + O(1/r^2)$$

Quasi Normal Modes (QNMs)



Scattering of Gravitational Radiation by a Schwarzschild Black-hole

THE discovery of pulsars and the general conviction that they are neutron stars resulting from gravitational collapse have strengthened the belief in the possible presence of Schwarzschild black-holes—or Schwarzschild horizons—in nature, the latter being the ultimate stage in the progressive spherical collapse of a massive star. The stability of these objects, which has been discussed in a recent report[1], ensures their continued existence after formation. Inasmuch as the infinite redshift associated with it and its behaviour as a one-way membrane make the

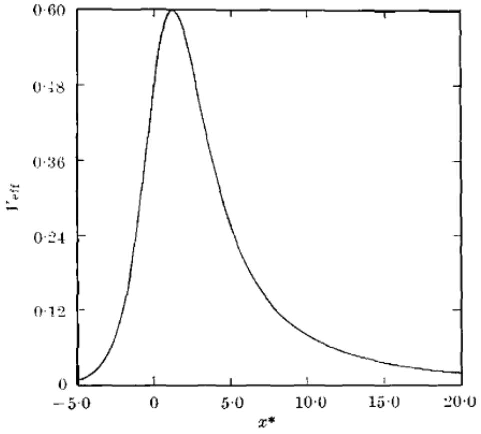
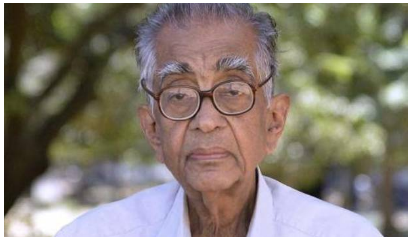


Fig. 1. The effective potential V_{eff} for the odd-parity gravitational waves of the lowest mode $l=2$ plotted against x^*.



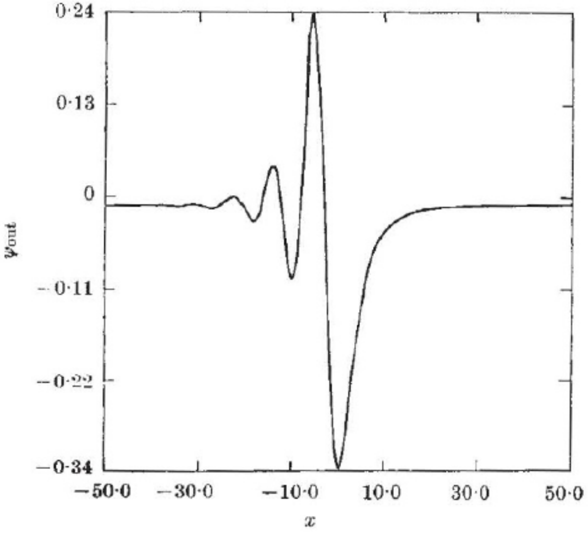


Fig. 3. The outgoing wave packet $\psi_{out}(x)$ at spatial infinity corresponding to the incident Gaussian wave packet $\psi_{in}(x) = e^{-ax^2}$ with $a=1$.

PHYSICAL REVIEW D VOLUME 5, NUMBER 12 15 JUNE 1972

Pulses of Gravitational Radiation of a Particle Falling Radially into a Schwarzschild Black Hole*

Marc Davis, Remo Ruffini, and Jayme Tiomno†
Joseph Henry Laboratories, Princeton University, Princeton, New Jersey 08540
(Received 20 December 1971)

Using the Regge-Wheeler-Zerilli formalism of fully relativistic linear perturbations in the Schwarzschild metric, we analyze the radiation of a particle of mass m falling into a Schwarzschild black hole of mass $M \gg m$. The detailed shape of the energy pulse and of the tide-producing components of the Riemann tensor at large distances from the source are given, as well as the angular distribution of the radiation. Finally, analysis of the energy going down the hole indicates the existence of a divergence; implications of this divergence as a testing ground of the approximation used are examined.

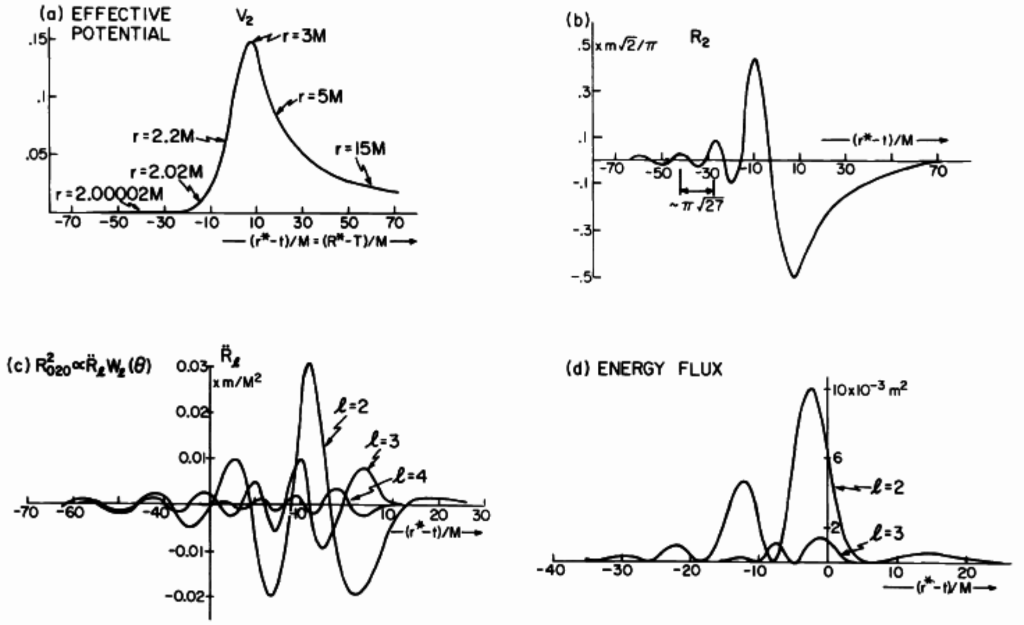


FIG. 1. Asymptotic behavior of the outgoing burst of gravitational radiation compared with the effective potential, as a function of the retarded time $(t-r^*)/M$. (a) Effective potential for $l=2$ in units of M^2 as a function of the retarded time $(t-r^*)/M=(T-R^*)/M$. For selected points the value of the Schwarzschild coordinate r is also given. (b) Radial dependence of the outgoing field $R_l(r,t)$ as a function of the retarded time for $l=2$. (c) $\ddot{R}_l(r^*,t)$ factors of the Riemann tensor components (see text) given as a function of the retarded time for $l=2,3,4$. (d) Energy flux integrated over angles for $l=2,3$; the contributions of higher l are negligible.

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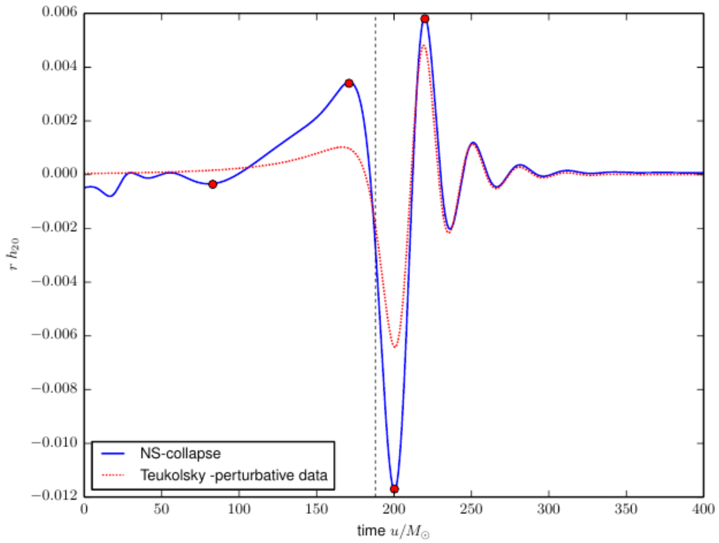
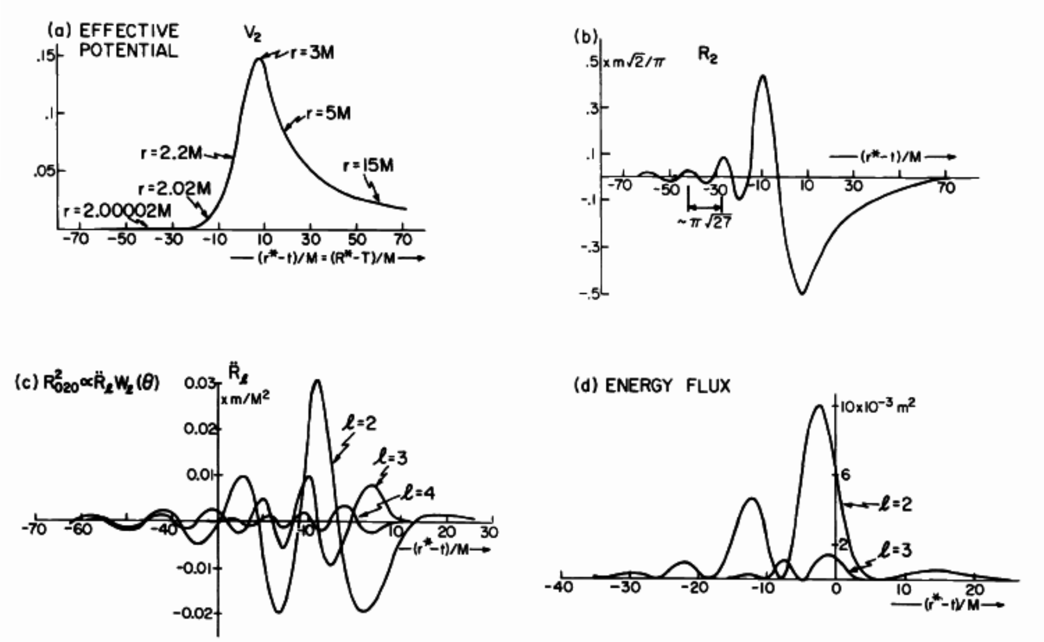


Signal from NS rotating collapse



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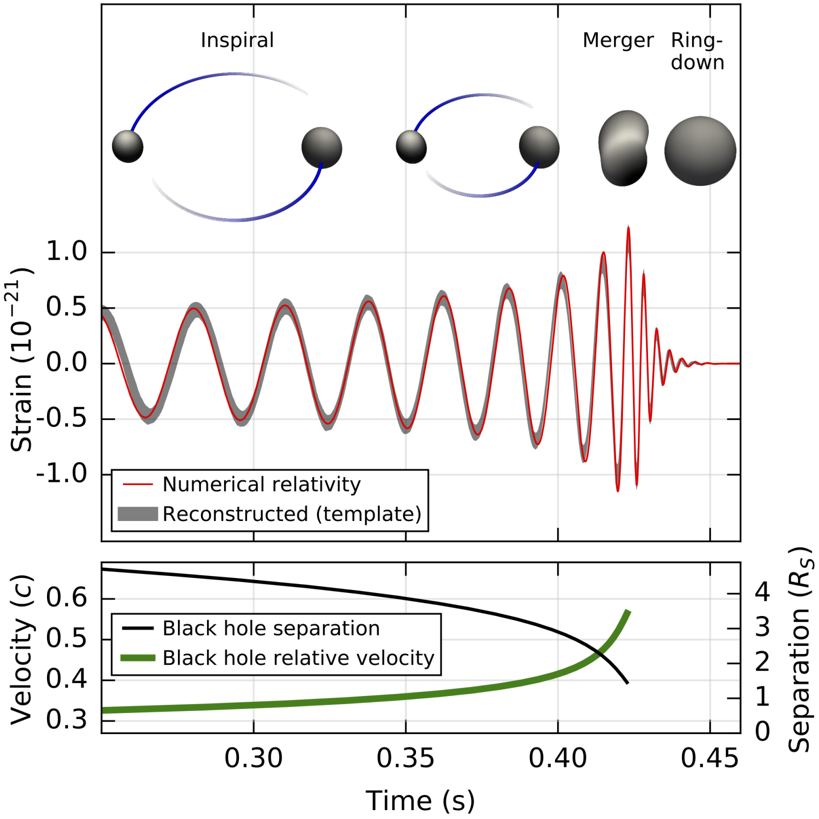
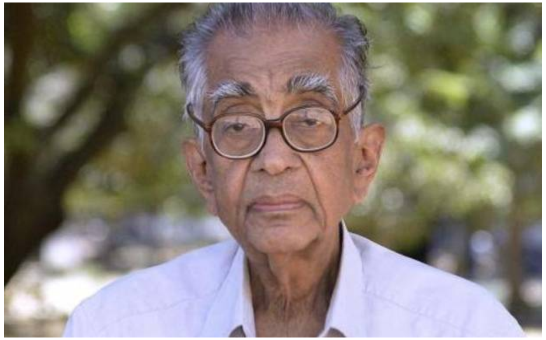
QNMs in binary black holes remnants





"It was a natural question then to ask: how does one see a black hole? So, using a computer, I scattered packets of gravitational waves from a black hole and the quasinormal modes emerged carrying the signatures of the black hole… this was theoretical. I had never dreamed that this would take on an aspect of reality some day,"

Prof. Vishveshwara (6 March 1938 – 16 January 2017)
[www.thehindu.com]

Origin of QNMs

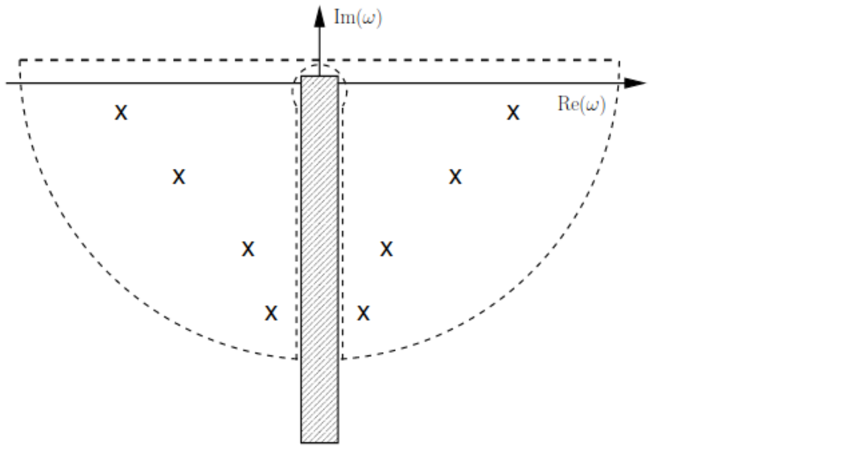
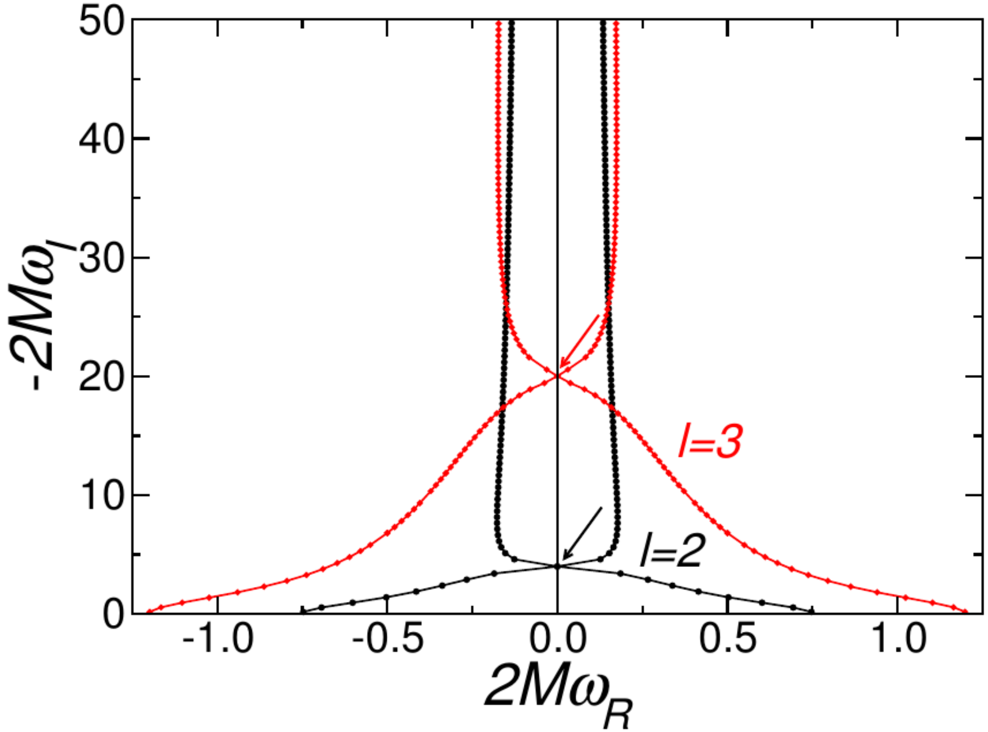


Figure 2. Integration contour for Eq. (36). The hatched area is the branch cut and crosses mark zeros of the Wronskian W (the QNM frequencies).



Review: https://arxiv.org/abs/0905.2975

Perturbations of spherical spacetimes

Consider the perturbation $h_{\alpha\beta}$ of a spherically symmetric background metric $g^{(0)}_{\alpha\beta}$ ($\mathcal{M} = M^2 \times S^2$) in some suitable coordinates (e.g. Schwarzschild):

$$g_{\alpha\beta} = g^{(0)}_{\alpha\beta} + h_{\alpha\beta} \tag{1}$$

Because of the background, the perturbation can be decomposed in scalar $Y_{\ell m}$, vectors $Z^{\ell m}_a$ and tensor $Z^{\ell m}_{ab}$ spherical harmonics with indexes (ℓ, m) and further separated between *even* (*electric-type*) and *odd* (*magnetic-type*) parity according to the behaviour under reflection trhough the origin: $(-1)^\ell$ and $(-1)^{\ell+1}$:

$$h_{\mu\nu} = h^{(e)}_{\mu\nu} + h^{(o)}_{\mu\nu} \tag{2}$$

For example, the decomposition of the even parity part reads ($A = 0, 1; a = 2, 3$)

$$h^{(e)}_{\mu\nu} = \begin{pmatrix} H_0 Y_{\ell m} & H_1 Y_{\ell m} & h^{(e)}_A Z^{\ell m}_a \\ '' & H_2 Y_{\ell m} & \\ '' & '' & r^2 K Y_{\ell m} \Omega_{ab} + r^2 G Z^{\ell m}_{ab} \end{pmatrix} \tag{3}$$

where Ω_{ab} is the metric on S^2 and the metric coefficients do not carry multipolar indexes for simplicity (A sum on (ℓ, m) is also understood).

Gauge invariant quantities (under infinitesimal coordinate transformations) can be identified from the above metric. Of particular importance are the two *scalar functions* for each multipole (suffix (ℓ, m) understood):

$$\Psi^{(e)}(t,r) \;\text{ and }\; \Psi^{(o)}(t,r) \tag{4}$$

The perturbed EFE lead to the *Regge-Wheleer-Zerilli (RWZ)* wave equation for the above scalar functions (one for each multipole (ℓ, m) that are all decoupled from each other):

$$\Psi_{tt} - \Psi_{xx} + V_\ell = S_{\ell m} \tag{5}$$

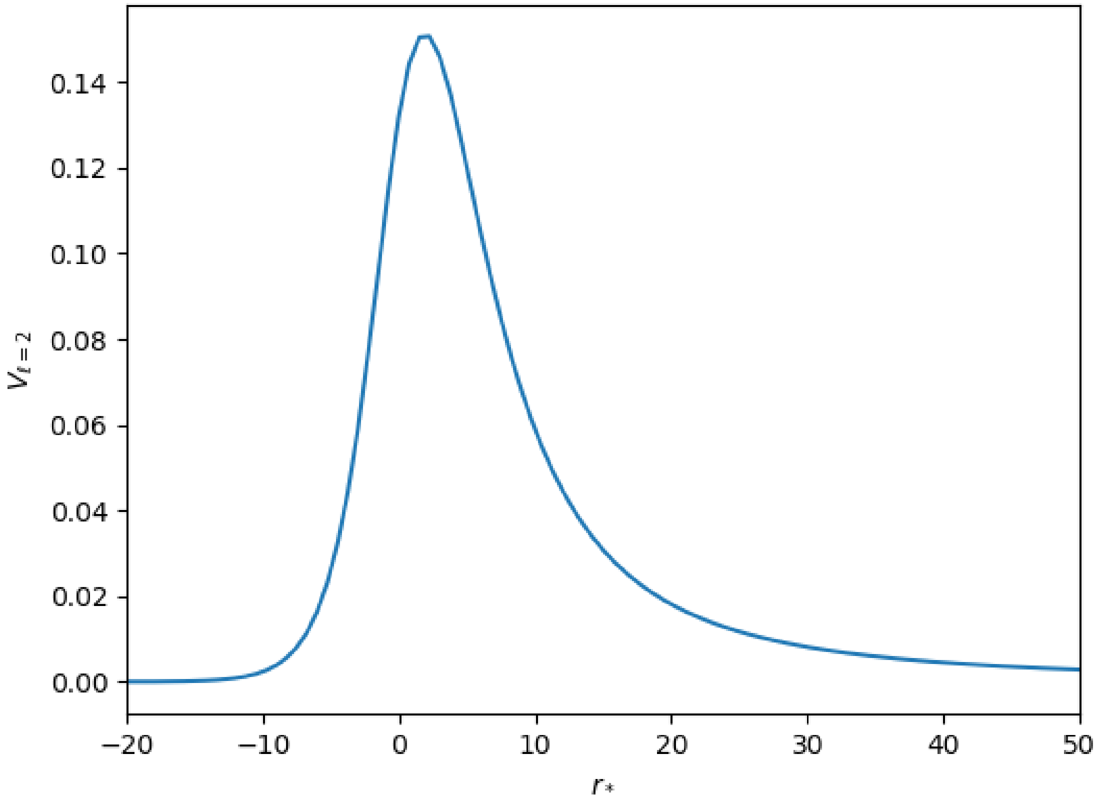
where x is the tortoise coordinate that maps $[2M, \infty)$ to $(-\infty, \infty)$

$$x = r + 2M \ln\left(\frac{r}{2M} - 1\right) \tag{6}$$

$S_{\ell m}$ is a source term from the stress-energy tensor, and V_ℓ is a potential determined by the background metric that for even and odd parity reads, respectively ($\Lambda := \ell(\ell+1)$)

$$V_\ell^{(e)} = A(r)\frac{\Lambda(\Lambda-2)^2 r^3 + 6(\Lambda-2)^2 M r^2 + 36(\Lambda-2)M^2 r + 72M^3}{r^3\left((\Lambda-2)r + 6M\right)^2} \tag{7}$$

$$V_\ell^{(o)} = A(r)\left(\frac{\Lambda}{r^2} - \frac{6M}{r^3}\right) \tag{8}$$



0.14
0.12
0.10
0.08
0.06
0.04
0.02
0.00
$V_{\ell=2}$
−20
−10
0
10
20
30
40
50
r_*

There is no dependence on m due to the spherical symmetry of the background. Among the linearized EFE, Eq.(5) plays a special role because its asymptotic solutions for large r represent the gravitational-wave degrees of freedom in the spin weighted spherical harmonics decomposition

$$h_+ - i h_\times = \frac{G}{c^4}\frac{1}{r}\sum_{\ell=2}\sum_{m=-\ell}^{\ell}\sqrt{\frac{(\ell+2)!}{(\ell-2)!}}\left(\Psi^{(e)}_{\ell m}(t) + i\,\Psi^{(o)}_{\ell m}(t)\right){}^{-2}Y_{\ell m}(\theta,\phi) \tag{9}$$

The RWZ problem in vacuum

The initial-boundary value problem with the RWZ requires to chose appropriate initial and boundary conditions. Because the RWZ potential tends to zero for both $x\to\pm\infty$ (horizon and spatial infinity), the asymptotic solutions at large (tortoise) radii are the solution of the "free" wave equation on the light cones.

By considering solutions with time dependence $\Psi\sim e^{-i\omega t}$ (or, equivalently, the Fourier modes), the RWZ equation can be cast in a form similar to the Schroedinger equation for stationary states,

$$\frac{d^2\tilde{\Psi}}{dx^2} + [\omega^2 - V_\ell]\tilde{\Psi} = 0 \tag{10}$$

However, since the RWZ potential is positive, no "bound states" can exists, and the spectrum is continuous. The physical requirement that no signals can come out from the the horizon, implies that the boundary condition at $x\to-\infty$ is an ingoing wave,

$$\tilde{\Psi}\sim e^{-i\omega x} \quad (x\to-\infty) \tag{11}$$

This boundary condition also follows from requiring smoothness. Requiring instead hat no signal can come in from spatial infinity, implies an outgoing wave for $x \to +\infty$

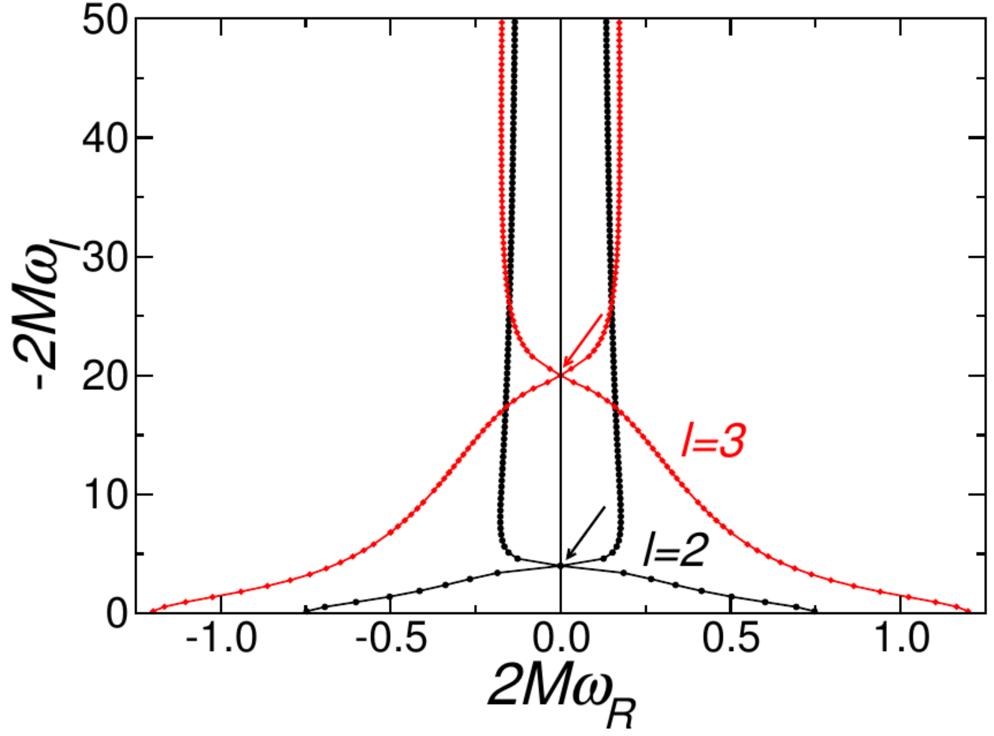
$$\tilde{\Psi} \sim e^{i\omega x} \quad (x \to +\infty) \tag{12}$$

With these boundary conditions, Eq.(10) admits solutions for a discrete infinity of *complex* frequencies ω_n with negative imaginary frequencies $\mathrm{Im}(\omega_n) < 0$. These damped modes are called *quasi-normal modes* (QNMs) appear also in other wave problems with *open* boundaries, and generically characterize *dissipative systems*. Differently from the normal modes of a vibrating string with "fixed" boundary conditions, QNMs do not form a complete set of eigenfunctions for the solution.

The presence of damped QNMs in the solutions suggests the stability of Schwatrzschild black holes under small perturbations (mode stability). The conclusion is correct, although the story is richer.

Some steps:

- Regge-Wheeler (1957) use a WKB analysis to argue that odd perturbations of the Schwarzschild spacetime are stable under the boundary conditions Eq.(11-12)
- Zerilli (1970) obtains the master equation for even parity (same WKB as above applies)
- Vishveshwara (1970) rules out perturbations growing in time because they diverge at the event horizon, if they fall at infinity.
- Chandrasekhar (1975) finds a map between the odd and even parity perturbations, and proves the QNMs are "isospectral"



50
40
30
20
10
0
-2Mω_I
-1.0
-0.5
0.0
0.5
1.0
2Mω_R
l=3
l=2

- Leaver (1986) formally identifies QNM as pole of a Green function
- Kay & Wald (1987) show that solutions with data of compact support are bounded
- Bachelot and Motet-Bachelot (1993) prove the existance of infinite number of QNMs

Solution by Laplace transform

The Cauchy problem specified by Eq.(5), boundary conditions (like Eq.(11-12)) and initial data $\Psi(0,x)=\psi(x)$ and $\Psi_t(0,x)=\psi_t(x)$ with compact support (or sufficiently localized) can be solved introducing the Laplace transform

$$\phi(s,x)=\int_0^\infty e^{-st}\Psi(t,x)\,dt \tag{13}$$

The Laplace transform is defined for positive, real $s>0$ and can be analytically continued into the positive complex plane. The equation for ϕ can be immediately found by integrating the RWZ,

$$\phi_{xx}-(s^2+V(x))\phi=F(s,x):=-s\psi(x)-\psi_t(x) \tag{14}$$

Two independent solutions $f_\pm(s,x)$ of the homogeneous equation ($F\equiv 0$) determine the unique Green function of the problem; the solution is

$$\phi(s,t)=\int_{-\infty}^{+\infty}G(s;x,x')F(s,x')dx'=\int_{-\infty}^{+\infty}\frac{f_-(s,x_-)f_+(s,x_+)}{W(s)}F(s,x')dx' \tag{15}$$

where $x_\pm = \frac{\max}{\min}(x, x')$ and $W(s)$ is the Wrongskian. The formal solution of the Cauchy problem is then obtained from the inverse Laplace transform

$$\Psi(t,x) = \frac{1}{2\pi i}\lim_{R\to\infty}\int_{\epsilon-iR}^{\epsilon+iR} e^{st}\,\phi(s,x)\,ds \tag{16}$$

where ϵ (real) is greater than the real part of all the singularities of ϕ .

The Laplace solution contains both the initial data (in $F(s,x)$) and the boundary conditions. The latter are implemented in the choice of the homogeneous solutions $f_\pm$. The integral in Eq.(16) can be performed using the residue theorem by chosing an appropriate contour in the complex plane, as determined by the analytical properties of ϕ .

The RWZ potential decays exponentially for $x \to -\infty$, it reaches a maximum and then decays as $1/x^2$ for $x \to \infty$. For the RWZ potential can be proven that (Bachelot&Motet-Bachelot 1993):

- f_- has poles only at negative real integers
- f_+ has a branch cut in the negative real axis due to the r^{-2} decay at large radii

The solution is then determined by different contributions

$$\Psi \sim \int_{\epsilon-iR}^{\epsilon+iR}(.) = \underbrace{\int_{\mathrm{large\,half-circle}}(.)}_{\mathrm{source\,term}} + \underbrace{\sum_k \mathrm{res}(.,s_k)}_{\mathrm{QNMs}} + \underbrace{\int_{\mathrm{branch\,cut}}(.)}_{\mathrm{Late-time\,tails}} \tag{17}$$

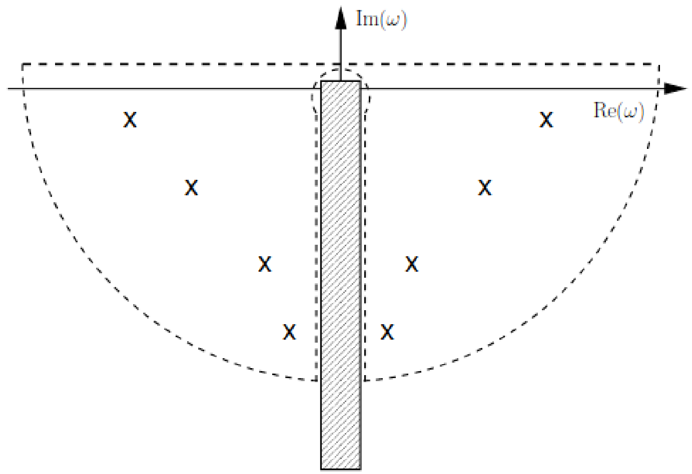


Figure 2. Integration contour for Eq. (36). The hatched area is the branch cut and crosses mark zeros of the Wronskian W (the QNM frequencies).

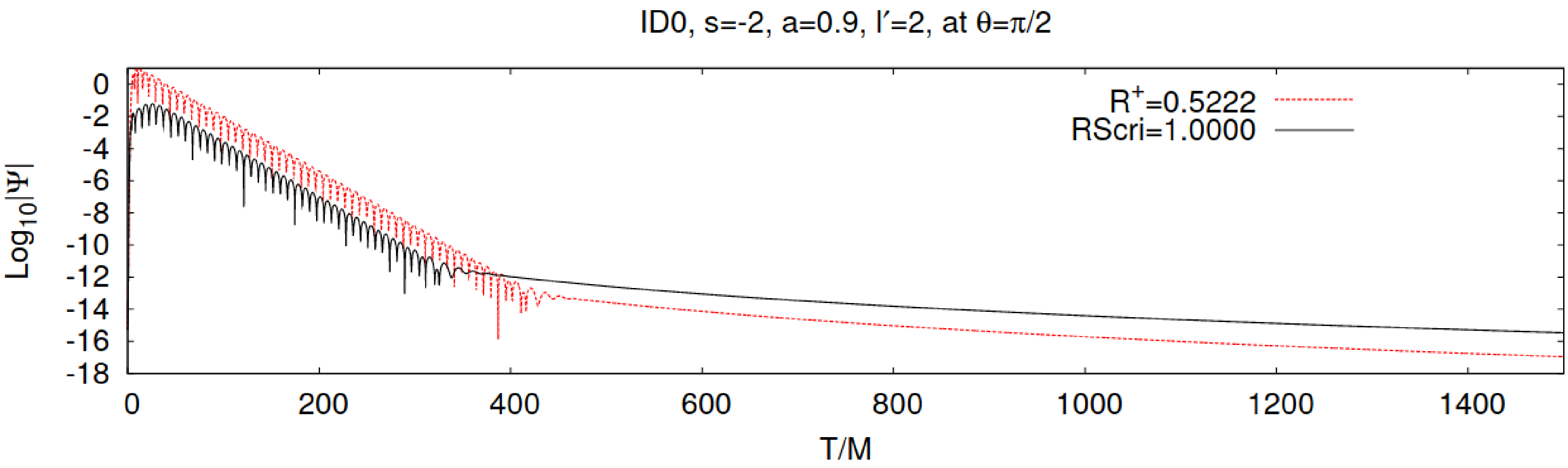


Figure 3. Evolution of the perturbation field at the horizon and $\mathscr{I}^+$ ($\theta = \pi/2$). The field is characterized by the quasi normal mode ringdown and a power law tail. The plot refers to a simulation of an axisymmetric gravitational perturbation ($s = -2$ and $m = 0$) with ID0, $l' = 2$ and $a = 0.9$.

Stability of black hole spacetimes

A fundamental question about exact stationary solutions of Einstein's field equations (EFE) like Minkowski, Schwarzschild or Kerr is their stability under small perturbations. Rigorous proofs are very nontrivial and usually built on several results.

A very rough scheme is the following:

- Linear mode stability: within linear perturbation theory one proves that the time evolution of each mode, say $\Psi_{\ell m}$, is bounded (in some norm) for a suitable class of initial data (say, with compact support).

- Linear stability: mode stability does not, in general, guarantee that a solution composed of an infinite sum of modes remains bound. Here one proves that all solutions to the linearised EFE remain bounded for all times by a suitable norm of their initial data. Mode stability is a necessary condition to linear stability.

- Nonlinear stability: here one considers the more general Cauchy problem in GR with initial data "near" Minkowski, Schwarzschild or Kerr, and shows that the solution remains bound.

The question of nonlinear stability of Kerr black holes is still open, although many positive results are available. An incomplete list is:

- [7] First argument for mode stability of Schwarzschild
- [5] Linear stability of scalar perturbation of Schwarzschild
- [8] Mode stability of Kerr
- [2] Nonlinear stability of Mikowski for asymptotically flat vacuum initial data
- [3] Linear stability of Schwarzschild
- [4] Linear stability of scalar perturbation of nonextremal Kerr BH
- [6] Nonlinear stability of Schwarzschild proven for a class of nontrivial perturbations
- [1] All extremal Kerr BH are unstable to gravitational perturbationalong their event horizon
- [?] Nonlinear stability of Schwarzschild

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[4] Mihalis Dafermos, Igor Rodnianski, and Yakov Shlapentokh-Rothman. Decay for solutions of the wave equation on Kerr exterior spacetimes III: The full subextremal case $|a| < M$. 2014.

[5] Bernard S. Kay and Robert M. Wald. Linear Stability of Schwarzschild Under Perturbations Which Are Nonvanishing on the Bifurcation Two Sphere. *Class. Quant. Grav.*, 4:893–898, 1987.

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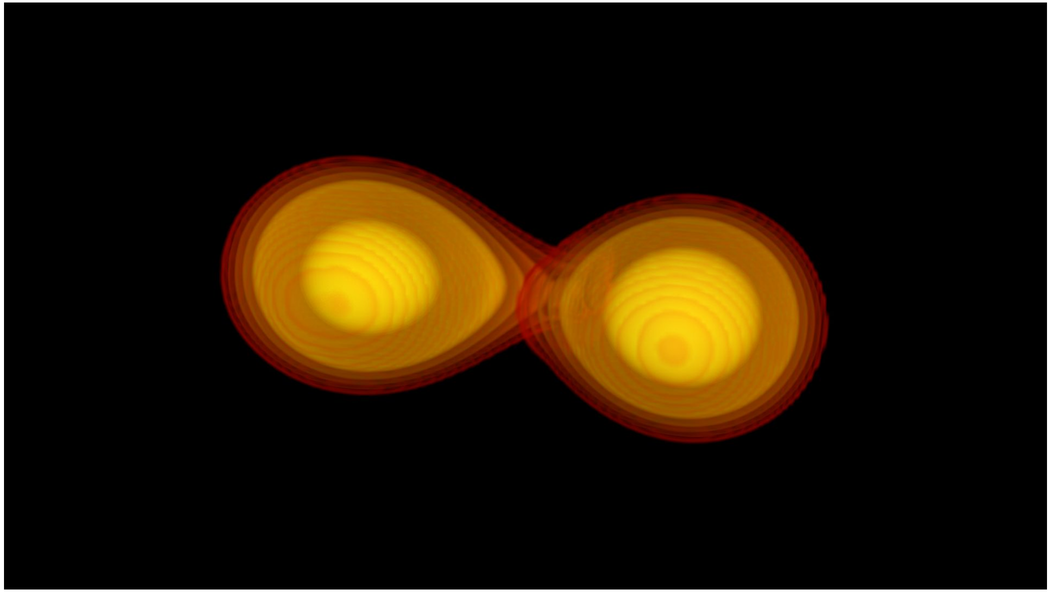
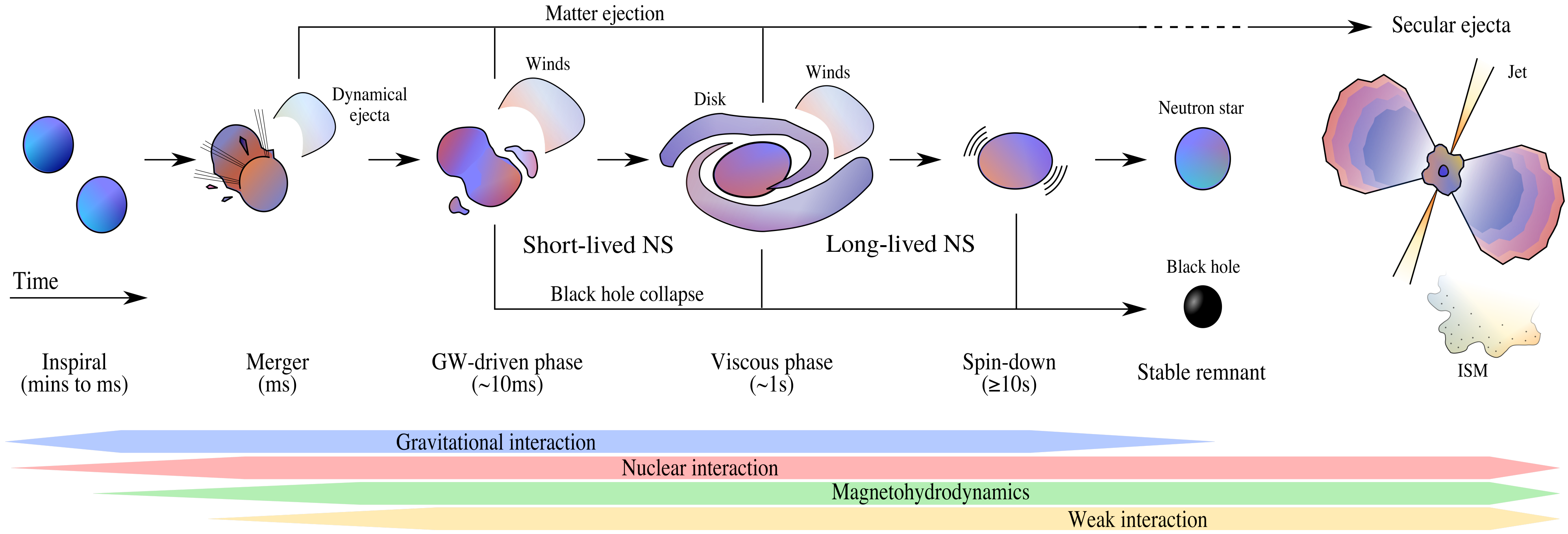
NS in binaries: tides



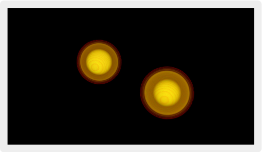
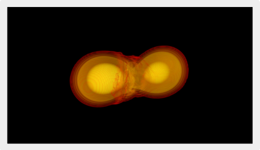
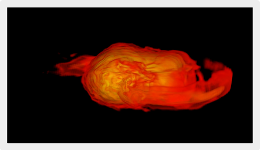
BNS mergers “(2-body dynamics)^4”

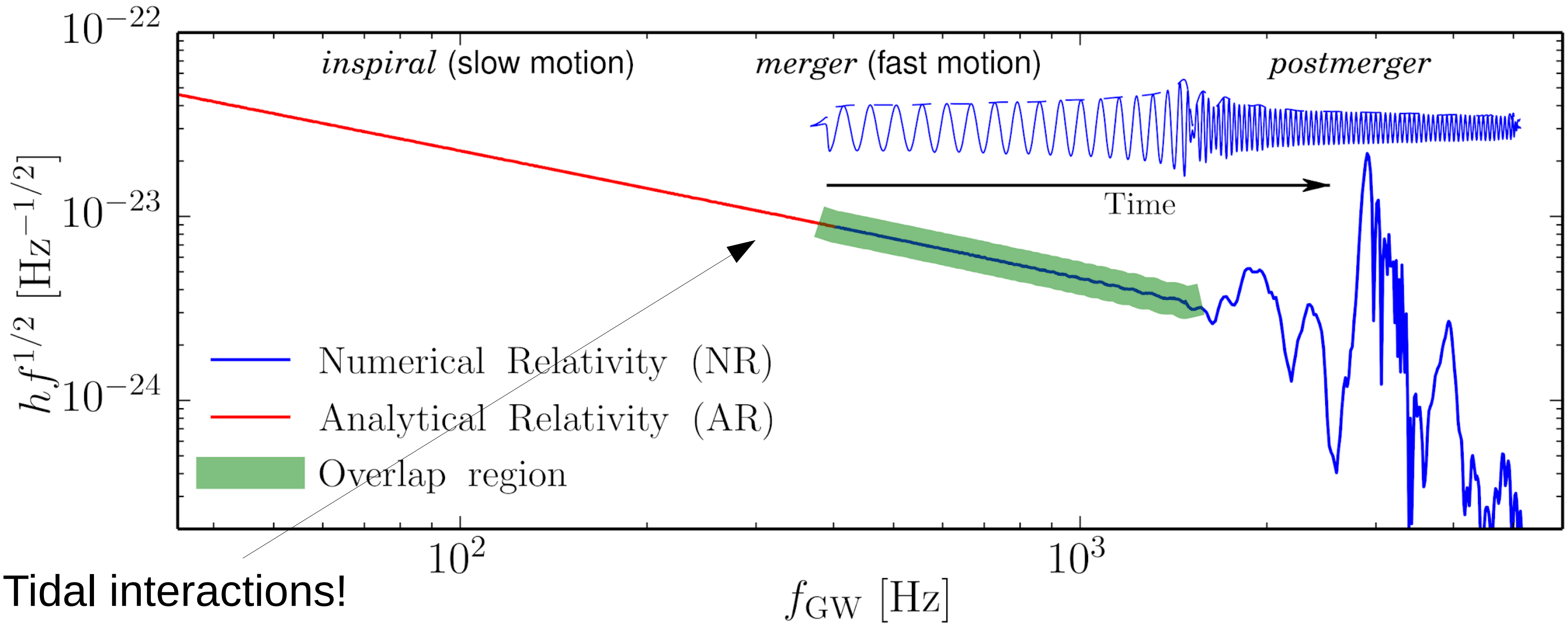


Annual Reviews of Nuclear and Particle Science
Radice, SB, Perego [https://arxiv.org/abs/2002.03863]

GERG Collection on Binary Neutron Star mergers
SB [https://arxiv.org/abs/2002.03863]

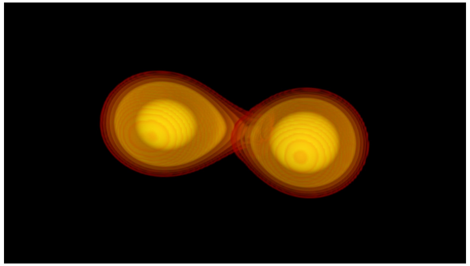
The gravitational-wave spectrum



Tidal interactions!

Tidal interactions in BNS



$$\kappa_2^T = \frac{3}{2}\left[\Lambda_A \left(\frac{M_A}{M}\right)^4 \frac{M_B}{M} + (A \leftrightarrow B)\right]$$

(Damour&Nagar 2009a 2009b)

Hamiltonian (Newtonian limit):

$$H_{\rm EOB} \approx Mc^2 + \frac{\mu}{2}\left[p^2 + A(r) - 1\right]$$

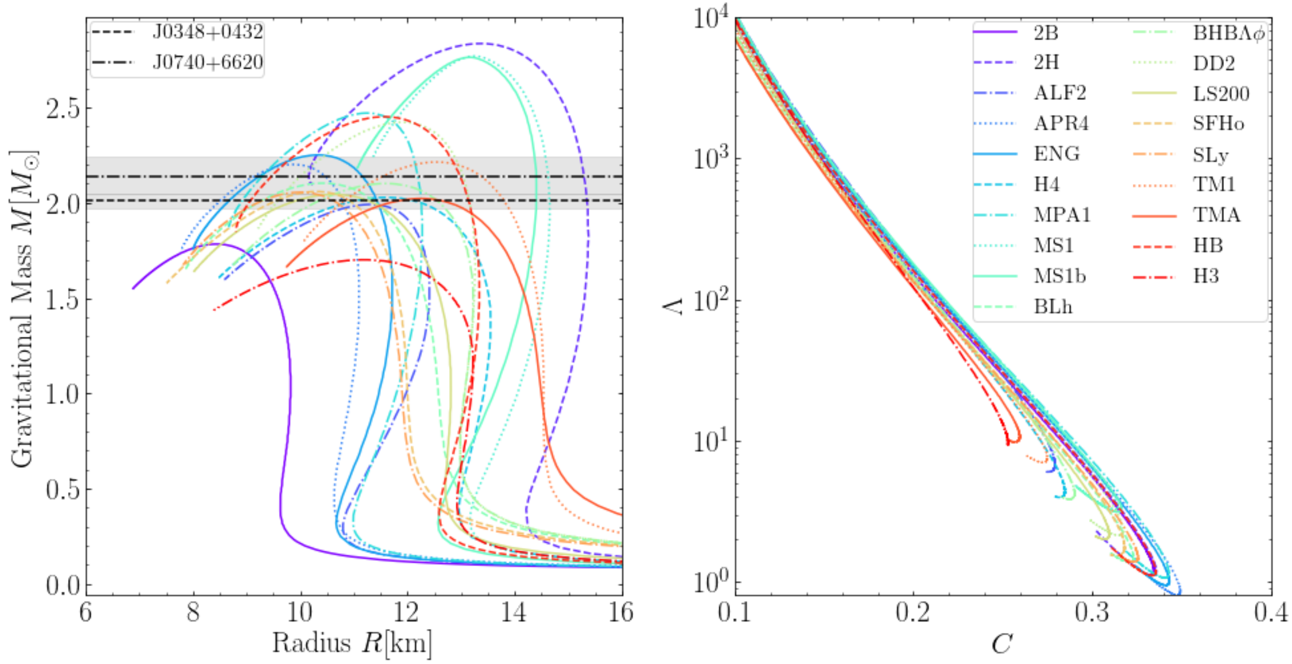
$$A(r) = 1 - \frac{2GM}{c^2 r} - \frac{\kappa_2^T}{r^6}$$

Waveform:

$$h \approx A f^{-7/6} e^{-i\Psi(f)} \approx A f^{-7/6} e^{-i\Psi_{\rm p.m.}(f) + i\frac{39}{4}\kappa_2^T (x(f))^{5/2}}$$

Key point: No other binary parameter (mass, radii, etc) enter separately the formalism at LO

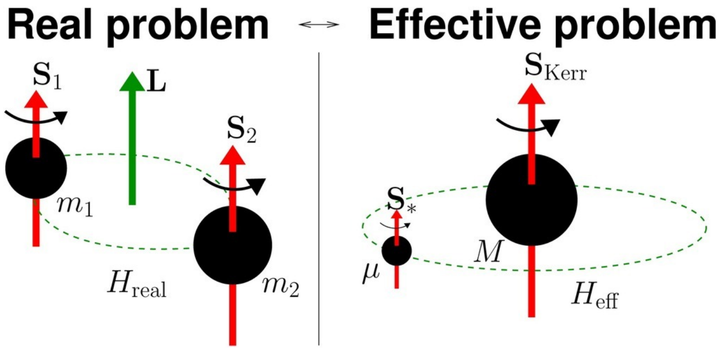
Tidal polarizability coefficients



[Hinderer 2007, Damour&Nagar 2009a, Binnington&Poisson 2009]

Effective-one-body framework in a nutshell

[Buonanno&Damour PRD 2000a, 2000b]

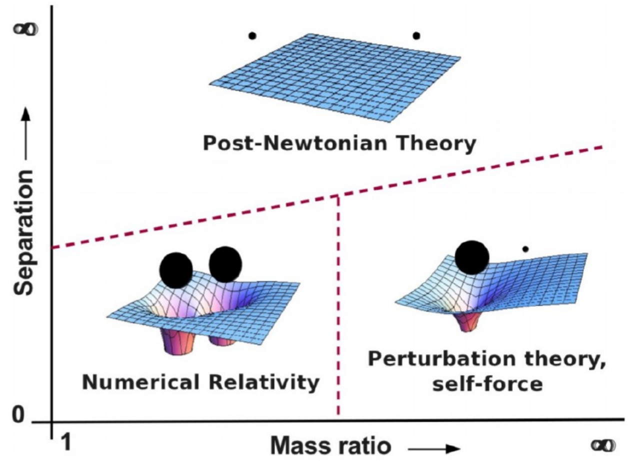


Credit: A.Taracchini/AEI

$$H_{\rm eff} \sim \mu\sqrt{A(u)(1+p_\phi^2 u^2)+p_{r*}^2}$$

$$A(u;\nu;\kappa_2^T) = A^0(u;\nu) + A^T(u;\nu;\kappa_2^T)$$

$$A^0(u;\nu) = 1 - 2u + \nu(\ldots$$



Credit: L.Barak

Factorized (resummed) PN waveform [Damour,Iyer,Nagar 2008]
Includes test-mass limit (i.e. particle on Schwarzschild)
Includes post-Newtonian and self-force results
Uses resummation techniques → predictive strong-field regime
Includes tidal interactions (→ BNS) [Damour&Nagar PRD 2010]
Flexible framework → NR informed

Compact binaries dynamics

The motion and radiation of a system of well separated, strongly self-gravitating ("compact"), bodies can be described by a "matching" approach which consists in splitting the problem into two (Damour 1983; Damour, Soffel, Xu 1991)

(i) the outer problem where one solves field equations in which the bodies are "skeletonized" by worldlines endowed with some global characteristics (such as mass, spin or higher-multipole moments)

(ii) the inner problem where one obtains the near-worldline behavior of the outer solution from a study of the influence of the other bodies on the structure of the fields in an inner world tube around each body

This approach can be used to obtain binary black hole dynamics in post-Newtonian (PN) formalism and to prove that the bodies' finite-size correction enters at 5PN.

Inner problem

Definition of multipolar tidal coefficients

Consider a static, spherically symmetric star of mass M perturbed by a stationary, external gravitational quadrupolar field $E_{ij} \sim \partial_i \partial_j \phi^{\text{external}}$. The star is expected to respond to the external field by developing a quadrupole moment Q_{ij}. This phenomeon is analogous to the electric polarizability of a medium that, placed in an external electric field, develops a dipole moment. Assuming a linear response, an (eletric-type) quadrupolar tidal coefficient is defined as

$$Q_{ij} = \mu_2 E_{ij} \quad (1)$$

A more general definition of μ_2 (valid also for other other multipoles) and the framework for the actual calculation can be obtained by the following argument.

In the star local frame and for large radii, the metric coefficient g_{00} (gravitational potential in the weak field) can be written as

$$\frac{1-g_{00}}{2} = -\frac{M}{r} + \frac{3}{2}\frac{Q_{ij}}{r^3}\left(n^i n^j - \frac{1}{3}\delta^{ij}\right) + \mathcal{O}\left(\frac{1}{r^4}\right) + \frac{1}{2}E_{ij}x^i x^j + \mathcal{O}(r^3) \tag{2}$$

The above expression shows that the tidal coefficient μ_2 in Eq.(1) can be obtained by matching the term growing as $\sim r^2$ to the term falling as $\sim 1/r^3$ of the asymptotics expression of the (perturbed) metric coefficient.

This procedure can be generalized. In the local frame of the body, define the *external gravito-electric and gravitomagnetic tidal moments*

$$G_L := G_{i_1 \dots i_\ell} \quad H_L := H_{i_1 \dots i_\ell} \tag{3}$$

as those *multipoles* of the perturbed metric that grow as r^ℓ. Similarly, the *internal mass and spin multipoles moments*

$$M_L \qquad S_L \tag{4}$$

are those that decay as $r^{-(\ell+1)}$. The multipolar tidal coefficients of the body are then postulated as those relating the linear response of the internal moments to the external ones

$$M_L = \mu_\ell G_L \qquad S_L = \sigma_\ell H_L \tag{5}$$

In a linearly perturbed, stationary star spacetime the asymptotics behaviour of the field uniquely defines these moments (Note: this is different from the vacuum case). In the following only the gravitoeletric sector is discussed since the magnetic sector is analogous.

Calculation of tidal Love numbers

Consider even parity, stationary perturbations of the TOV metric $g^{(0)}_{\alpha\beta}$

$$g_{\alpha\beta} = g^{(0)}_{\alpha\beta} + h^{(e)}_{\alpha\beta} \tag{6}$$

The $h^{(e)}_{00}$ coefficient of the perturbed metric can be expressed in terms of a function H that is directly related to the logarithm of the enthalpy perturbation. The perturbative equation for H is

$$H'' + C_1 H' + C_0^{(\ell)} H = 0 \tag{7}$$

with

$$\begin{aligned} C_0^{(\ell)} &= e^{2\lambda}\left[\frac{\ell(\ell+1)}{r^2} + 4\pi\,(\rho+P)\frac{d\rho}{dP} + 4\pi\,(5\rho+9P)\right] - 4\phi'^2 \\ C_1 &= \frac{2}{r} + e^{2\lambda}\left[\frac{2m}{r} + 4\pi r\,(p-\rho\,)\right] \end{aligned}$$

(A similar equation hold for odd parity perturbations). In the star interior, Eq.(7) needs to be solved numerically together with the background equations and by specifying a EOS. In the star exterior, $\rho = P = 0$ and $m = M$, and the equation reduces to the *associated Legendre equation* with variable $x = r/M - 1$. The general solution can be expressed in terms of the *associated Legendre functions*

$$H^{\rm outer} = a_P \hat{P}_{\ell 2}(x) + a_Q \hat{Q}_{\ell 2}(x) \tag{8}$$

The coefficients a_P and a_Q are to be determined by the boundary conditions, in particular by the matching with the interior solution. The ratio $a_\ell := a_Q/a_P$ can be determined by requiring the continuity of the logarithmic derivative at the surface

$$y_\ell(r) = \frac{r H'(r)}{H(r)} \tag{9}$$

i.e.

$$y_\ell^{\rm inner}(R) = y_\ell^{\rm outer}(R) = (1 + x_R) \frac{\hat{P}'_{\ell 2}(x_R) + a_\ell \hat{P}_{\ell 2}(x_R)}{\hat{Q}'_{\ell 2}(x_R) + a_\ell \hat{Q}_{\ell 2}(x_R)} \tag{10}$$

with $x_R = R/M - 1 = 1/C - 1$. Note this is a nontrivial statement to check, since it depends on the EOS (and the regularity of the matter fields at the surface, e.g. the sound speed) and on the fact that the perturbed star surface does not coincide with the background star radius. Solving Eq.(10) for a_ℓ gives

$$a_\ell = -\frac{\hat{P}'_{\ell 2}(x_R) + C y_\ell(R) \hat{P}_{\ell 2}(x_R)}{\hat{Q}'_{\ell 2}(x_R) + C y_\ell(R) \hat{Q}_{\ell 2}(x_R)} \tag{11}$$

This coefficient can be now directly related to the tidal coefficient μ_ℓ. The asymptotic behaviour of the outer solution is determined by

$$\hat{P}_{\ell 2}(x) \sim \left(\frac{r}{M}\right)^{\ell+1} \qquad \hat{Q}_{\ell 2} \sim \left(\frac{M}{r}\right)^{\ell+1} \tag{12}$$

such that the growing and falling part of the perturbation are

$$\left(h_{00}^{(e)}\right)^{\text{growing}} \sim a_P \left(\frac{r}{M}\right)^{\ell+1} Y_{\ell m} \qquad \left(h_{00}^{(e)}\right)^{\text{falling}} \sim a_P \left(\frac{r}{M}\right)^{-(\ell+1)} Y_{\ell m} \tag{13}$$

The matching gives (reintroducing the constants G and c)

$$(2\ell-1)!!\, G\, \mu_\ell = a_\ell \left(\frac{G\,M}{c^2}\right)^{2\ell+1} \tag{14}$$

$G\,\mu_\ell$ has dimension of $[\text{length}]^{2\ell+1}$. The *tidal Love numbers* are defined as the dimensionless combination

$$k_\ell := \frac{1}{2} a_\ell\, C^{2\ell+1} = -\frac{1}{2} C^{2\ell+1}\, \frac{\hat{P}'_{\ell 2}(R/M-1) - C y_\ell(R) \hat{P}_{\ell 2}(R/M-1)}{\hat{Q}'_{\ell 2}(R/M-1) + C y_\ell(R) \hat{Q}_{\ell 2}(R/M-1)} \tag{15}$$

The *tidal polarizability parameters* of a star often employed in gravitational-wave astronomy are defined as

$$\Lambda_\ell := \frac{2\,k_\ell}{(2\ell-1)!!C^{2\ell+1}} \tag{16}$$

Outer problem

Effective action

Up to 5PN order ($\mathcal{O}(v/c)^{10}$) the motion of two body compact bodies of mass M^A $A=1,2$ is described by the effective action

$$S = \int \frac{R}{16\pi G} - \sum_{A=1}^{2} \int M^A d\tau_A \tag{17}$$

where an "apportune regularization" must be introduced to deal with the point-mass source term in the EFE. Note that the calculation of the 5PN dynamics is not yet completed: for nonspinning bodies, the conservative dynamics is fully known at 4PN, while the waveform at 3.5PN.

Finite-size effects enter at 5PN and the action needs to be augment with the term

$$S_{\text{nonminimal}} = \sum_{A,\ell} \left[\frac{1}{2}\frac{\mu_\ell}{\ell!}\int (G_L^A)^2\, d\tau_A + \frac{1}{2}\frac{\ell}{(\ell+1)}\frac{\sigma_\ell^A}{\ell!}\frac{1}{c}\int (H_L^A)^2 d\tau_A \right] \tag{18}$$

Tidal Lagrangian at leading order

At leading order the tidal Lagrangian for body A is given by

$$L_T^A \sim \mu_2 \, (G_L^A)^2 \tag{19}$$

where the external tidal moment is calculated on the worldline z_A^a of the body as

$$G_L^A \sim \partial_L U^{\text{external}} = \partial_L \left(\frac{G\, M^B}{|r_{\text{AB}}|} \right) \tag{20}$$

with $r_{\text{AB}} = |z_A^a - z_B^a|$. The calculation uses the formalism of symmetric trace-free (STF) tensors for multipolar expansions, and in particular the expression

$$\partial_L^A \left(\frac{1}{r_{\text{AB}}} \right) = (-1)^\ell (2\ell - 1)!! \, \frac{\hat{n}^L}{r_{\text{AB}}^{\ell+1}} \tag{21}$$

where $\hat{n}^L$ is the STF projection of the of the unit vectors $n^a = (z_A^a - z_B^a) / r_{\text{AB}}$. The result is

$$L_T^A \sim + \sum_\ell \frac{(2\ell - 1)!!}{2} \mu_A \frac{(G\, M^B)^2}{r_{\text{AB}}^{2\ell+2}} = + \sum_\ell k_\ell^A G\, (M^B)^2 \frac{R_A^{2\ell+1}}{r_{\text{AB}}^{2\ell+2}} \tag{22}$$

The interaction is proportional to the Love numbers (or to the tidal polarizability parameters Λ_ℓ^A), it is attractive, and it is *short-range*, e.g. the first term scales as $\sim 1/r^6$.

Was this useful? Quick self-check:

- What are the two main characteristics of compact objects?
- What is compactness? Make a table with order of magnitude values for mass, radius, compactness, average density for WD, NS, BH (without books/googling)
- What is the origin of pressure support in WD? What is the order of magnitude of the critical density? What are the values of the adiabatic index above/below the critical density?
- What is the physical origin of the Chandrasekhar mass? Can you provide an order of magnitude argument for the existance of $M_{\rm Ch}$?
- What is the Buchdal limit?
- What is the maximum NS mass $M_{\rm max}$? Can you give an upper bound?
- What is the difference between the WD and the NS EOS?
- What is a black hole? Why and How black holes form from stars?
- What is the LSO? And how can you estimate the merger frequency of black hole binaries?
- What is the RWZ equation? What boundary conditions are usually imposed?
- Are black holes stable? How can one formulate the stability problem?
- What are the Love numbers and the tidalpolarizability parameters? How can they be computed?

Now you can do the proposed exercises and check the references mentioned in the lectures!