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What do we know about neutron stars?

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Abstract. The Galaxy contains some 10^9 neutron stars which are thought to be born in supernova explosions of massive stars, and which we observe as radio pulsars if isolated and as x-ray binaries if attached to a close companion. Young radio pulsars are also observed as γ-ray pulsars, even as ultrahard γ-ray pulsars, and occasionally as infrared, optical, and x-ray pulsars. Binary neutron stars are probably spun down by the wind of their companion until an accretion disc forms which confines the corotating magnetosphere. One then sees an x-ray source which may flicker via the ejection of two relativistic electron-positron beams, and pulse if the spin period is sufficiently long and/or the x-ray photosphere is non-axisymmetric *w.r.t.* the spin axis. This emission by a fast-rotating magnet may be similar to the one by active galactic nuclei. Old neutron stars can also emit x-ray bursts, and γ-ray bursts.

The present overview discusses this interpretation and alternative ones, together with the main theoretical tools. It outlines the arguments by which one derives a neutron star mass of $(1.4 \pm 0.2)\, M_\odot$, radius of $10^{1\pm0.2}$ km, surface magnetic field of $6 \cdot 10^{12\pm0.5}$ G and pulsar age of several 10^6 yr. Even in exotic sources like Sco X-1, SS 433 and Cyg X-1, the presence of a neutron star may well be indicated.

Key words: pulsars—binary neutron stars—supernova remnants—pulsing x-ray sources

1. Introduction

Neutron stars are an important stellar population, both in number and in energy output. They contribute significantly to the galactic radio, x-ray, γ-ray and ultrahard ($\lesssim 10^{16}$ eV) radiation and may even be responsible for the cosmic rays, *i.e.* for the extremely relativistic charged particles with energies upto 10^{20} eV, as a consequence of the enormous free energy stored in their rotation at formation. Their birth is (always?) marked by a gigantic supernova explosion; and even when they are 'dead', *i.e.* sufficiently spun down, they can still emit powerful x-rays and (bursts of) γ-rays through mass accretion from a nearby companion, or dense cloud.

In spite of the enormous amount of data collected on pulsars, and on the pulsing binary x-ray sources, in the past 15 years, there are almost as many different theoretical models for their explanation as there are independent workers; and several of the fundamental mechanisms are still controversial. Are the winds of neutron stars (mainly) composed of hydrogen or of electrons and positrons; are their magnetic fields frozen in, regenerated or decaying; do their moments of inertia fluctuate? How are the electrons accelerated to Lorentz factors between 10^2 and 10^9 (inferred from their synchrotron and inverse-Compton spectra)? Where exactly and how is their radiation produced? Why do binary x-ray sources flicker?

Unique answers to these questions would not only help in reducing the relevant literature, but could at the same time serve as a guide to related astrophysical problems such as the quasar phenomenon (of active galactic nuclei), the phenomenon of bipolar flows in star-forming regions, and perhaps also that of biconical planetary nebulae. In all these situations, we observe pairs of antipodal jets emerging from an almost pointlike central source and pushing their way to extreme distances through the ambient intergalactic or interstellar medium, with (at least partially) nonthermal spectra, considerable linear polarization and remarkable stability (*cf.* Kundt 1984, 1982a). If Sco X-1 and SS 433 are prototypes of such quasar-like behaviour on a stellar scale, and if they are fast-spinning binary neutron stars, we may speculate that the quasar phenomenon relies on a fast-rotating magnet as the central engine. In the four mentioned cases of an AGN, bipolar flow, planetary nebula and binary neutron star, this rotating magnet would be realized respectively by the magnetized (super-) massive core of a massive disc, by a young magnetic white dwarf, and by a young (magnetic) neutron star. If it turns out to be true that pulsar winds are similar in composition to the jets of quasars and newly-forming stars, we learn that the mechanism of the central engine must work under a large variety of different astrophysical conditions, and therefore must not depend on very special configurations. This—admittedly speculative—generalization underscores the need for a thorough and simple understanding of the pulsar phenomenon.

There are a number of excellent reviews on neutron stars in the literature. For instance, their stellar structure has been summarized by Källman (1979) and by Baym & Pethick (1979); pulsar theory can be found in Sutherland (1979) and in Michel (1982), see also the books by Manchester & Taylor (1977), Melrose (1980), and Shapiro & Teukolsky (1983), further Radhakrishnan (1982) and the symposium edited by Sieber & Wielebinski (1981); binary x-ray sources are treated in the books by Sanford *et al.* (1982) and Lewin & van den Heuvel (1983), in particular in the contributions to the latter by Lewin & Joss (Galactic bulge sources), Rappaport & Joss (x-ray pulsars) and van den Heuvel (binary evolution), see also Börner (1980) and Joss & Rappaport (1984). In the present overview I do not include any historical background, nor do I endeavour any bibliographical completeness. Instead, my main emphasis will be on the achieved results and on the most important ambiguities in their interpretation, *i.e.* on the explicit or tacit assumptions and conclusions made by various authors. In such a young and fundamental field of astrophysics as that of neutron stars, we are often in danger of dismissing the proper explanation by adhering too tightly to conventional wisdom; I shall mention all the viable alternatives I am aware of.

2. Neutron star birth

A neutron star forms when the degenerate core of a massive star collapses, owing essentially to its own weight. The estimated critical mass M_{crit} (of the progenitor star at formation) above which this happens has grown throughout the past ten years; its present value lies between 6 and 12 $M_\odot$; *cf.* van den Heuvel (1983). A star of initial mass $M < M_{crit}$ develops into a white dwarf, or suffers explosive disruption. The latter possibility—explosive disruption—is offered by stellar-evolutionary calculations and seems to be needed to explain the many 'empty' supernova shells. On the other hand, if we dismiss many supernova explosions as the birth sites of neutron stars, we run into the dilemma that there are more neutron stars in the sky than birth events. Note that a neutron star may be shielded by the (weak but heavy) wind of a companion star as long as its spin is high enough to impede accretion; this so-called dormant or grindstone stage will be discussed below. The companion star, in turn, may be hidden in or behind the cloud from which the double star formed. Alternatively, note that a pulsar whose spindown power ($\sim \Omega\dot{\Omega}$) at birth is some 10^2 times lower than that of the Crab can be hidden in the radio noise of its envelopping supernova shell (Radhakrishnan 1982). It is therefore misleading to call a shell 'empty' without specifying on what criteria and at what sensitivity limit this statement holds. For instance, Cas A may well harbour another binary pulsar (Kundt 1980).

2.1. Supernova explosions

When the core of a (sufficiently) massive star collapses, it can liberate its gravitational binding energy on the free-fall timescale (of order ms). This binding energy is the difference between gravitational and all other forms of energy except rest-mass energy :

$$E_{bind} = E_{grav} - (E_{int} + E_{rot} + E_{turb} + E_{mag}),$$

with $E_{rest} = Mc^2 = 10^{54.5}$ erg $(M/1.4M_\odot)$, $E_{grav} = (0.2 \pm 0.1)\, Mc^2$ (obtained within General Relatively Theory, measured as redshift) and $E_{rot} = I\Omega^2/2 \lesssim 10^{52.5}$ erg, where I = moment of inertia $\lesssim 10^{45}$ g cm^2, $\Omega \leqslant \Omega_{crit} = (GM/R^3)^{1/2} \approx 10^4$ s^{-1}. The ratio E_{int}/E_{grav} equals 1/2 for cold degenerate matter within Newtonian theory, and unity for cold, relativistically degenerate matter. The net static binding energy of a (cold, magnetised) neutron star is therefore larger than its maximal E_{rot} plus E_{turb}, and smaller than $E_{grav}/2$ (at Newtonian approximation), hence $E_{bind} \approx 10^{53.2\pm0.5}$ erg.

Most pulsars seem to be born with angular velocities $\Omega_* \leqslant 3\cdot10^2$ s^{-1} which are much slower than the critical angular velocity $\Omega_{crit} = 10^4$ s^{-1} for centrifugal support (exceptions are the recently discovered two fastest pulsars, with $\Omega = 4\cdot10^3$ s^{-1}

and 10^3 s^{-1}, which show however exceptionally small spindown). Such (comparatively) slow rotators should have liberated most of their static binding energy in the form of gravitational and/or neutrino radiation. Less easily radiated is the rotational energy, because of the angular momentum constraint. If the collapsing core passes through extreme rotation, the corresponding free energy amounts to $E_{rot} = I\Omega^2/2 = 5 \cdot 10^{52}$ erg I_{45} ($I_{45} := I/10^{45}$ g cm^2 = moment of intertia in units of 10^{45} g cm^2). When this energy is transferred to the collapsing outer shell of the star, via the strained toroidal magnetic field—a magnetic spring—this shell will fly off explosively without any further energy input (Bisnovatyi-Kogan *et al.* 1976; Kundt 1976).

This suggestion of a magnetically driven supernova explosion has been widely ignored in the literature, perhaps because it is still beyond the reach of the fastest computers unless truncated to two dimensions. Plausible as it may sound, its effectiveness depends on a few constraints : in particular the spring must couple the core to a comparable moment of inertia of the surrounding shell in order to achieve a transfer without bounce; and the spring must break, *i.e.* the field must recombine, soon after takeoff of the accelerated shell.

Tapping the rotational energy of the collapsing core would not work if a nuclear-chemical explosion transferred enough radial momentum to the outer shell before the magnetic torque reaches its critical value for angular momentum transfer. In this case, the driving piston would be neutrinos (*cf.* Hillebrandt 1984). Most workers in the field pursue this path. Both pistons, the magnetic and the neutrino piston, are Rayleigh-Taylor unstable : denser layers are pushed by underlying low-density ones. Yet there is a significant difference between the two : magnetic Rayleigh-Taylor instabilities tend to produce one-dimensional patterns, say filaments, whereas pure density instabilities tend to produce isotropic patterns, like explosion clouds, or volcanic plumes. If there were two types of supernova mechanisms, their remnant shells would be expected to show different morphologies. To me, all supernova shells look filamentary, but in many cases we have to await higher resolution. Note that the efficiency of the magnetic piston drops to zero when the mass of the ejected shell—or rather its moment of inertia—falls below that of the remaining core, as it would for collapsing white dwarfs with an outer shell of mass $\leqslant 0.2\, M_\odot$.

Supernovae tend to eject some to several $M_\odot$ at velocities of $10^{4\pm0.3}$ km s^{-1}, corresponding to kinetic energies of $10^{51\pm1}$ erg. (The kinetic energy of one $M_\odot$ ejected at 10^4 km s^{-1} amounts to 10^{51} erg.) These numbers apply both to spectral studies of young supernovae and to the historical ($\leqslant 10^3$ yr old) supernova shells for which one knows both their age and present radius. In comparison, the newborn Crab pulsar had a rotational energy of only $3 \cdot 10^{49}$ erg (Kundt & Krotscheck 1980).

Some years after their explosion, supernovae can become strong radio emitters for several years; see Weiler *et al.* (1983). The necessary relativistic electrons may be supplied by a central neutron star, and squeeze their way out through the expanding filamentary shell.

2.2. *Supernova shells*

Several hundred years after their explosion, supernovae will reappear as expanding shells, or arcs of filamentary shells. In well observed cases like the Crab and Cas A, the 3-d velocities of individual bright spots have been measured and found to conform with strictly radial expansion, *cf.* figure 1. The shell of Cas A shows even three different velocity systems : optical 'knots' with v_{exp} between $3 \cdot 10^3$ and 10^4 km s^{-1}, optical 'flocculi' with 10 times smaller velocities, and radio brightness peaks with intermediate velocities (Tuffs 1983). Such a multicomponent structure is reminiscent of filamentary 'bullets' moving almost unimpeded through the ambient ISM, the lower-velocity systems being wake phenomena. They are not adequately described by the one-component Sedov-Taylor detonation wave, nor are the different observed temperature regimes (radio, optical, x-ray) consistent with a Sedov wave.

For these reasons, a 2-component model has been proposed by Kundt (1983, 1985b); see also Kundt & Gopal-Krishna (1984). It predicts an exponential slowdown (rather than Sedov's $t^{-3/5}$ power law), with an e^{-1}-folding time of $3 \cdot 10^3$ yr $\rho_{-25}^{-1/2}$ (where ρ is the average ambient mass density). Accordingly, supernova shells 'die' suddenly, at a typical age of 10^4 yr, and much earlier ($\sim 10^3$ yr) if located inside a dense galactic cloud like perhaps Tycho, Kepler and MSH 15-52. This modification of our understanding, if confirmed, implies that we have often overestimated the ages of supernova

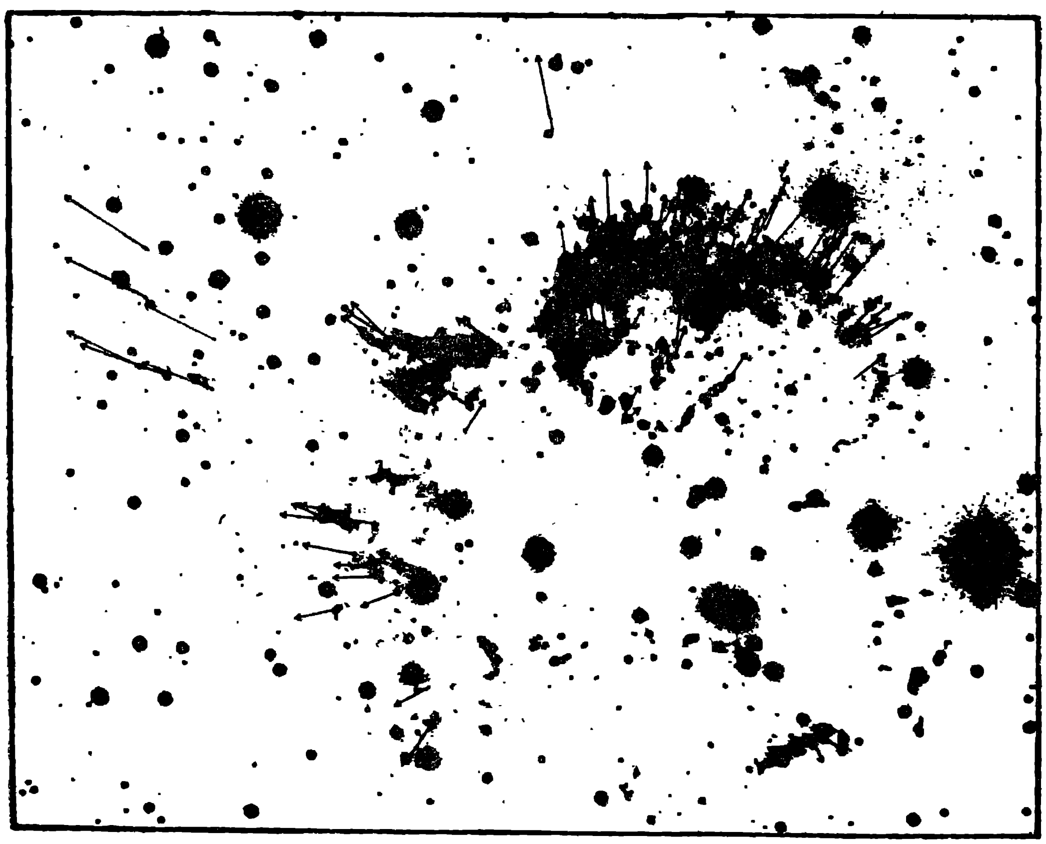

Figure 1. Optical image of the supernova shell Cas A, with the proper motion vectors (within 50 yr) of 117 fast-moving knots drawn in (copied from Kamper & van den Bergh 1976). The velocity vectors demonstrate an ordered radial motion in projection; the equally measured radial velocities show that also the 3-d velocities are radial, whereby the main shell is decelerated by 6% *w.r.t.* the foreriding knots in the north-east. The emission shell has a radius of 5*l* yr (*d*/2.8 kpc).

shells, and hence underestimated their birthrate (see also Srinivasan & Dwarakanath 1982). Note that shell remnants are found around all (four) pulsars of age $\lesssim 10^4$ yr.

The 2-component supernova shell model implies a low sweeping efficiency ($\approx 10^{-2.5}$) and a correspondingly larger shell radius of near-free expansion. It removes the seeming age discrepancy between the (second youngest) pulsar PSR 1509–58 and its supernova shell G 292.0 + 1.8 (= MSH 15–52); *cf.* Katz (1983). It also offers a straight-forward explanation for the fact that many supernova shells have a bright and a dark hemisphere : If the shells inherit the orbital velocity of their progenitor star, they can have non-negligible center-of-mass velocities ($\approx 10^2$ km s^{-1}) compared with their expansion velocity, and ram pressure heating ($\sim \rho v^3$) must be considerably stronger on the side 'facing the (interstellar) wind'.

2.3. *Statistics*

If the birth of a neutron star liberates some 10^{52} erg in the form of kinetic and electromagnetic energy, our telescopes should be able to see it unless it happens inside or behind a molecular cloud. This argument applies independently of the mass of its progenitor star. If neutron stars can form from white dwarfs, as is often suggested (Isern *et al.* 1984), this enormous explosion energy would imply excessive velocities ($v \gtrsim 3 \cdot 10^4$ km s^{-1}) of the escaping stellar remnants which have not been observed. For this reason and the ones given above, I question the possibility that white dwarfs contribute significantly to supernova explosions; for an independent reasoning see van den Heuvel (1983).

The birthrate of neutron stars should therefore be smaller than or equal to the rate of supernova explosions, and to the birth rate of supernova shells. The birthrate $\dot{N}_{PSR}$ of pulsars depends on the unknown beaming factor f which has often been assumed unreasonably small ($f = 0.2$, *cf.* Narayan & Vivekanand 1983). Vivekanand & Narayan (1981) find $\dot{N}_{PSR} = 1/(20 \pm 5)$ yr for the Galaxy; see also Kundt (1977) where a fan beam of banana-shaped profile is suggested, and Jones (1980).

But massive stars, the progenitors of neutron stars, are rarely (or never?) single. This fact is not only implied by the huge peculiar velocities of pulsars, runaway stars, and (probably) supernova shells, but also brought out by direct observations (Stone 1979) and explained by stellar formation theory. In particular, the so-called 'single' massive stars are quite often runaways with unexplained residuals in their radial velocities, erratic light curves, and excess nonthermal radio emission (above 6 cm wavelength : Abbott *et al.* 1984); they may well have a neutron star companion. Double or multiple stars can form from their progenitor disc when the latter fractionates, due to the bar mode instability (Bodenheimer 1981).

The birthrate of neutron stars will therefore be roughly twice as high as the birth rate of pulsars : $\dot{N}_{n*} \approx 1/10$ yr, and we expect at least as many supernova events. This is more or less in accord with Tammann's (1977) statistics of supernovae but hardly in accord with supernova shell statistics (Srinivasan & Dwarakanath 1981). Quite likely we are only aware of less than 50% of all supernova shells; deep radio surveys and special sources (Sco X-1 : Kundt & Gopal-Krishna 1984) point in this direction. Note that this case of the 'missing shell remnants' would be worsened if we dismissed 'empty' shells as the birthsites of neutron stars.

3. Pulsar data

Pulsars are observed as scintillating (pointlike) radio sources emitting regular pulses at repetition periods P between 1.5 ms and 4.3 s, with smeared-out radio luminosities (near 400 MHz) between 10^{25} and 10^{31} erg s^{-1}. More than 340 pulsars are currently known. Most of them have low-frequency cutoffs in their spectra near $\nu_l = 10^2$ MHz (Izvekova *et al.* 1981), the Crab being a notable exception (with $\nu_l < 10$ MHz). Optical and γ-ray pulses have so far only been seen from the Crab and the Vela pulsar (γ : Graver & Schönfelder 1983), but ultrahard γ-ray pulses (10^{11} eV$<E<10^{17}$ eV) were received from the Crab, Vela, and the nearby pulsar 0950 + 08 [Gupta *et al.* (1978); see figure 6].

The pulsars form a disc population of the Galaxy, with a scale height of some 350 pc and a possible clustering in a ringlike domain of radius $\lesssim 5$ kpc (Lyne 1982). One therefore observes an isotropic distribution in the sky out to some 200 pc from the solar system which then changes into a Milky-way distribution beyond, but gets increasingly incomplete beyond distances of 500 pc. The proper motions of 26 well-observed pulsars have an inferred 3-d mean value of $\langle (\mathbf{v})^2 \rangle^{1/2} = 2 \cdot 10^2$ km s^{-1} (Lyne *et al.* 1982), much larger than that of any other stellar population.

The temperatures of pulsars are much lower than predicted by cooling calculations of the 1970s : No neutron star has so far been discovered as a thermal x-ray source (above 0.2 keV), corresponding to surface temperatures below $2.5 \cdot 10^6$ K for the Crab and $\ll 10^6$ K for all the other 72 supernova remnants and pulsars searched, 10 of which were within 500 pc distance of the sun (Helfand 1983; Helfand & Becker 1984).

3.1. *Pulsar magnetic moments*

If a pulsar loses its rotational energy $E_{\rm rot} = I\Omega^2/2$ in accord with the (rotating, magnetic) dipole radiation formula

$$-\dot{E}_{\rm rot} = (2/3c^3)\, \ddot{\mu}_\perp^2 = 2\,\Omega^4\, \mu_\perp^2/3c^3, \qquad ...(1)$$

where $\mu_\perp := \mu \sin \psi$ (with $\mu := BR^3$, $\psi := \sphericalangle(\mathbf{\Omega}, \boldsymbol{\mu})$) is the transverse magnetic dipole moment and a dot denotes time derivative, then its period $P = 2\pi/\Omega$ should lengthen with time according to $P^2 - P_*^2 \sim t$, for $\mu_\perp =$ const, and its surface magnetic field B should be given by

$$B_\perp = (3\,Ic^3\, P\dot{P}/8\pi^2\, R^6)^{1/2} = 10^{12}\, G\, (P\dot{P}/10^{-15}\, s)^{1/2}. \qquad ...(2)$$

Correspondingly, for $P \gg P_*$, a pulsar's age τ should be given by $\tau = P/2\dot{P}$, and its braking index $n := \Omega\ddot{\Omega}/\dot{\Omega}^2$ should be equal to 3. For young pulsars ($\tau < 10^6$ yr), the spindown age (or 'characteristic' age) τ does indeed seem to be a good measure of age, but not so for $\tau \gg 10^6$ yr; nor are $B_\perp$ and n statistically independent of τ.

An apparent way to generalize the above evolution is to assume an exponentially decaying transverse magnetic dipole moment $\mu_\perp$ (Flowers & Ruderman 1977). If $\mu_\perp$ decays in proportion to $\exp(-t/2\tau_\mu)$ then we find instead

$$P/P_* = 1 + (\tau_\mu/\tau_*)\,(1 - \exp(-\,t/\tau_\mu)), \qquad \ldots(3)$$

and

$$\tau/\tau_* = \exp(t/\tau_\mu)\,[1 + (\tau_\mu/\tau_*)\,(1 - \exp(-t/\tau_\mu))], \qquad \ldots(4)$$

so that the period P freezes out at a time $t \gtrsim \tau_\mu$, and τ grows exponentially with t. Of course, equations (3) and (4) do not distinguish between a decaying μ and a decaying $\sin\psi$.

What is known theoretically about the temporal behaviour of a neutron star's magnetic field? As a first approach, textbooks sometimes consider the magnetic fields of a main-sequence star frozen during its collapse and obtain the correct order-of-magnitude of a pulsar's field (Zeldovich *et al.* 1983). Ruderman & Sutherland (1973) refined this approach by estimating the self-generated field of the degenerate core of a neutron star's progenitor, and Levy & Rose (1974) pointed out that such a dynamo would be expected to have a significant dipole moment. More recently, one can also find statements that a neutron stars's field may be generated during the collapse that leads to its formation (van den Heuvel & Taam 1984), or that a neutron star's field may be thermally generated after its formation (Blandford *et al.* 1983). The last suggestion is certainly interesting. But if a surface field of some 10^{13} G is to be generated within 10^3 yr from the heat current of a neutron star of surface temperature $\lesssim 3\cdot10^6$ K—corresponding to the Crab pulsar—the conversion efficiency η of heat into magnetic energy would have to be of order $\eta \approx 10^{-2} B_{13}^{2}/T_{6.5}^{4}\, t_{10.5}$ Such a high efficiency compares unfavourably with those of the solar and the planetary dynamos which are likely below 10^{-4}, whereby the planetary dynamos may be driven by sedimentation rather than by cooling (Moffatt 1978). A magnetic dynamo could alternatively be driven by differential rotation (spindown). Below we shall encounter phenomenological arguments that strengthen the constant-field hypothesis.

3.2. *Pulsar ages*

In the subsequent discussion of the data I shall ignore the possibility of field (re-) generation, and try to obtain a satisfactory fit to equations (3) and (4). In figure 2, the published values of $(P\dot{P})^{1/2}$ are plotted against those of $P/2\dot{P}$ for almost 300 pulsars. Straight lines of fixed $P\dot{P}$, P and $\dot{P}/P^3$ are drawn through the extreme points of the distribution where three short-period pulsars were ignored. The lines suggest that young pulsars ($\tau \lesssim 10^5$ yr) have $B_\perp$ = const whereas old pulsars ($\tau \gtrsim 10^6$ yr) evolve according to P = const. The broken line ($\dot{P}/P^3$ = const) is a line of constant spindown power; pulsars above it would not be detected. There may likewise be an observational discrimination against very short periods P; and many of the $\gtrsim 30$ pulsars for which $\dot{P}$ has not yet been published may have large τ.

Nevertheless, the following conclusions can be drawn from figure 2 : (i) Young pulsars have remarkably uniform magnetic surface fields, between $3\cdot 10^{12}$ and $2\cdot 10^{13}$ G; (ii) if individual pulsars evolve according to the law (3), the transition age τ_μ from $B_\perp$ = const to P = const must fall between $5\cdot 10^4$ yr and $5\cdot 10^5$ yr; (iii) along a strip $|P - P_\infty| \leqslant$ const, the number density of pulsars decreases roughly as $\tau^{-1/2}$ with spindown age, not as τ^0 as demanded by equation (4).

Consequently, if $\mu_\perp$ decays, it must do so on average within less than $3\cdot 10^5$ yr. This conclusion is at variance with the one drawn from the histogram of $B_\perp$-values, figure 3, which suggests a decay age of $0.8\cdot 10^7$ yr (if a uniform population of pulsars is assumed : Lyne 1982, figure 10).

A third and again different age estimate for 'magnetic field decay' has been obtained from a comparison of spindown ages with 'kinetic' ages $t_{\rm kin} = |z - z_*|/|\dot{z}|$, where z is the height above the galactic disc, and z_* is the (unknown) height at birth which can be larger than 10^2 pc, as evidenced by the Crab and by PSR 0611 + 22. The data show that $t_{\rm kin}$ saturates near 10^7 yr whereas τ increases beyond 10^9 yr. From

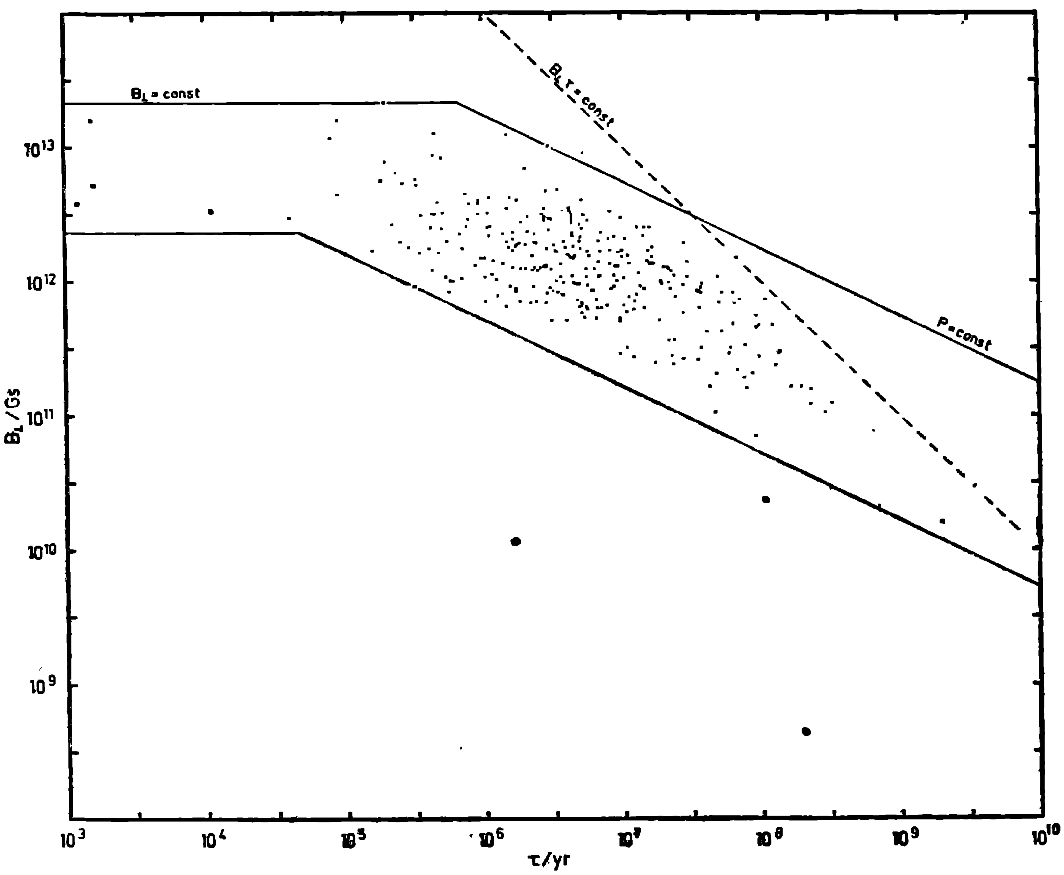

Figure 2. Values of $(P\dot{P})^{1/2} \sim B_\perp/B_*$ *versus* spindown age $P/2\dot{P} =: \tau$ for some 300 pulsars taken from Manchester & Taylor (1981), except for five recently discovered ones. $B_* = (3Ic^3\, 10^{-15}s/8\pi^2R^6)^{1/2} = 10^{12}$G $(I/R^6)_9{}^{1/2}$ according to equation (2). Almost all points fall within a strip bordered by $B_\perp$ = const and P = const respectively, and further restricted by the detectability constraint $B_\perp\tau \lesssim 10^{20}$G s.

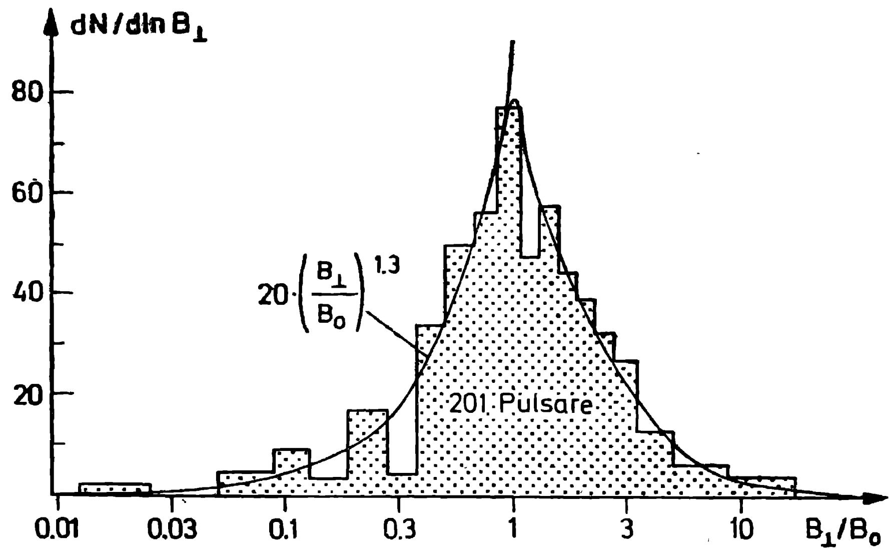

Figure 3. Histogram of transverse dipole moments $\mu_\perp$, or equivalent field strengths $B_\perp$, with $B_0 = B_*$ as in figure 2. Note that this histogram can be obtained from figure 2 by projecting the distribution parallel to the τ-axis; it is much wider for the whole population than for restricted (spindown) age groups.

their 26 data points, Lyne *et al.* (1982) conclude at a 'field decay time' of $5 \cdot 10^6$ yr, which is inconsistent with the two preceding determinations. Their conclusion loses its cogency, however, when allowance is made for larger values of $| z_* |$, and for the presence of two populations of pulsars : the younger and older 'brothers' in a (massive) binary system (Kundt 1977). At the same time, Anderson & Lyne's (1983) seeming correlation between pulsar proper motions and transverse magnetic dipole moments at birth $\mu_\perp$ disappears when the 'magnetic decay time' τ_μ is chosen of order $3 \cdot 10^5$ yr rather than equal to $8 \cdot 10^6$ yr.

To sum up, the kinetic estimate shows that pulsar ages do not exceed some 10^7 yr, so that τ must be an overestimate of true age (for $\tau > 10^6$ yr). Quite likely, $\mu_\perp$ decays. But a single (exponential) field decay is inconsistent with figure 2 because as stated under (iii), the density of points drops as $\tau^{-1/2}$ instead of being constant: pulsars disappear on their way to large τ, for some reason different from field decay (Kundt 1981a). Perhaps pulsars align, *i.e.* turn into parallel rotators. Further evidence against magnetic field decay (within less than some 10^8 yr) comes from the —often much older—x-ray pulsars (Kundt 1981a) and from the γ-ray bursters (Lipunov *et al.* 1982) both of which tend to have inferred surface field strengths above 10^{12} G; *see also* Kundt *et al.* (1985).

3.3. Period irregularities

Our equation (3) above yields a satisfactory description of pulse arrival times for most pulsars most of the time, with an accuracy of prediction that can reach $0.6 \cdot 10^{-12}$ (in the case of PSR 1133 + 16 : Gullahorn & Rankin 1982), or even $3 \cdot 10^{-14}$ (in the case of the 1.56 ms PSR 1937 + 214). Yet most or all pulsars show 'noise' in their pulse arrival times when one monitors them for sufficiently long.

Such noise can manifest itself in the form of discrete jumps ('glitches') in the pulse frequency, of size $\Delta\Omega/\Omega \leqslant 2\cdot 10^{-6}$, and/or jumps in its time derivative, of size $\Delta\dot{\Omega}/\dot{\Omega} \lessapprox 3\cdot 10^{-2}$, or both; whereby so far, all glitches in excess of $|\Delta\Omega|/\Omega > 3\cdot 10^{-9}$ have been positive whereas smaller jumps have shown either sign of $\Delta\Omega$. On top of the glitches, the pulse arrival times show quasi-continuous noise. This noise may well be a superposition of many unresolved microglitches, and may dominate the power of irregularities. Table 1 contains, among others, all pulsars for which glitches $\Delta\Omega/\Omega$ in excess of 10^{-7} have been observed.

Noise in the pulse arrival times means that either (i) the moment of inertia varies, or (ii) the braking torque is unsteady, or (iii) different components of the neutron star rotate at different angular velocities, with varying coupling strength. Which of the three is the dominating cause?

When the first glitches were observed, a plausible explanation seemed to be quakes in the solid crust, or core, of the slowing rotator which approaches the (spherical) shape of smallest moment of inertia. But an approach to the state of strongest gravitational binding involves a liberation of the binding energy, probably via damped oscillations, whose dissipated heat should be 'visible' in the form of x-rays. No such excess heat has been observed, in particular from the nearby Vela pulsar with the largest and most frequent glitches (six within 12 years). Moreover, our understanding of neutron star matter has changed in the sense that no solid core (quantum

Table 1. Some data for some fast and/or young and/or irregular pulsars, in order of increasing pulse period. The first column gives the trivial name, (X or IR standing for pulsed x-rays or infrared); $\tau := P/2\dot{P}$; $B_\perp$ is defined in equation (2); $n := \Omega\ddot{\Omega}/\dot{\Omega}^2$; $\Delta\Omega/\Omega$ and $\Delta\dot{\Omega}/\dot{\Omega}$ describe observed jumps; Δt estimates the repetition time of large glitches; and P_{orb} gives the orbital period of binary pulsars. The pulsing sources named *X* are not strictly 'pulsars'; they draw their power from accretion rather than rotation.

Name	Position	P/s	τ/yr	$B_\perp$/G	n	$\Delta\Omega/\Omega$	$\Delta\dot{\Omega}/\dot{\Omega}$	Δt/yr	P_{orb}/d
ms	1937+214	0.00156	$2\cdot 10^8$	$4\cdot 10^8$					
	1953+29	0.00613	$\geqslant 2\cdot 10^6$	$\leqslant 0.6\cdot 10^{11}$					~120
Crab	0531+21	0.0331	10^3	$4\cdot 10^{12}$	2.52	$2\cdot 10^{-8\pm 0.3}$	$2\cdot 10^{-4}$	6	
LMC	0540−69.3	0.0503	$2\cdot 10^3$	$5\cdot 10^{12}$					
binary	1913+16	0.0590	10^8	$2\cdot 10^{10}$					0.32
X	0538−66	0.069	—	—					16.66
	0530−67	0.07							
Vela	0833−35	0.0892	10^4	$3\cdot 10^{12}$	$0.6\cdot 10^2$	$2\cdot 10^{-6}$	$\left\{\begin{matrix}4\\0.7\end{matrix}\right\}\cdot 10^{-2}$	2.4	
	1509−58	0.150	$2\cdot 10^3$	$2\cdot 10^{13}$	2.83				
	0655+64	0.196	$2\cdot 10^9$	$2\cdot 10^{10}$					1.03
X+IR		0.286							
	1641−45	0.455	$4\cdot 10^5$	$3\cdot 10^{12}$		$2\cdot 10^{-7}$	$2\cdot 10^{-3}$	>2	
	1325−43	0.533	$3\cdot 10^5$	10^{12}		$1.2\cdot 10^{-7}$			
	2224+65	0.683	10^6	$3\cdot 10^{12}$		$1.7\cdot 10^{-6}$	$<0.6\cdot 10^{-3}$		
	1508+55	0.740	$2\cdot 10^6$	$2\cdot 10^{12}$	$3\cdot 10^3$	$2\cdot 10^{-10}$	$-0.6\cdot 10^{-2}$		
	0820+02	0.865	10^8	$3\cdot 10^{11}$					1232

*The data are taken from Manchester & Taylor (1981), Gullahorn & Rankin (1982: irregularities), Backus *et al.* (1982 : irregularities), Lohsen (1981 : Crab), Downs (1981 : Vela), Mc Culloch *et al.* (1983 : 6th glitch of Vela), Backer *et al.* (1913 : ms pulsar), Ashworth *et al.* (1983 : ms pulsar), Boriakoff *et al.* (1983 : 6 ms pulsar), Manchester *et al.* (1985: 1509–58), IAU telegrams, and conference reports.

crystal) is expected to exist (Baym & Pethick 1979), and that the crust has a different elastic structure from laboratory bodies by being much less compressible than shearable, whereby the density dictates the (Fermi) pressure, (Krotscheck *et al.* 1975). Equally unlikely are sudden enormous δ-shaped external torques (exerted *e.g.* by infalling bodies) that would almost periodically spin up (not spin down) the Vela and other pulsars.

There remains the possibility first advocated by M. Ruderman and collaborators and more recently elaborated by Alpar *et al.* (1982, 1984a, b, c) that the neutron star contains superfluid components whose spindown lags behind that of the rest of the star due to vortex creep. Such "weakly" and "superweakly pinned" shells suddenly share their angular momentum with the rest of the star during a glitch, but decouple again thereafter so that $\dot{\Omega}$ increases during the glitch. Alpar *et al.* (1984c) argue that the superfluid in the core of a neutron star is strongly coupled to its charged component via electron-magnetic-vortex scattering whereas the (pinned) superfluid in the inner crust is not. Its moment of inertia I_p therefore satisfies

$$I_p/I = \Delta\dot{\Omega}/\dot{\Omega} \leqslant 1 \qquad \text{...(5)}$$

and ranges between 10^{-4} and $4 \cdot 10^{-2}$ for different pulsars and events, *cf.* table 1. The maximum (relative) size of a glitch is similarly given by I_u/I times the repetition time Δt in units of (four times) the spindown time $\tau = -\Omega/2\dot{\Omega}$:

$$\Delta\Omega/\Omega \approx (I_u/I)\,(\Delta t/4\tau), \qquad \text{...(6)}$$

where I_u is the moment of inertia of a component which is uncoupled between two successive glitches. For the Vela, $\Delta\Omega/\Omega = 2 \cdot 10^{-6}$ and $\Delta t/4\tau = 0.6 \cdot 10^{-4}$ imply $I_u/I = 3 \cdot 10^{-2}$.

The glitches are the more spectacular but perhaps less important part of timing irregularities. Figure 4 shows the braking index $n := \Omega\ddot{\Omega}/\dot{\Omega}^2$ of 20 pulsars plotted against their spindown age, where positive values are represented by triangles 'pointing' upward. Surprisingly one finds an almost equi-distribution of positive and negative values which is bounded by the relation

$$|\,n\,| \lesssim \tau/3 \cdot 10^2 \text{ yr}. \qquad \text{...(7)}$$

Why does n differ from its dipole value 3 ?

The points in figure 4 were obtained after three to 10 years of monitoring and may well 'go down' for longer monitoring intervals. If they remained fixed, the equation $n = \pm\,\tau/\tau_c$ with $\tau_c = 3 \cdot 10^2$ yr would integrate to

$$\dot{\Omega}/\dot{\Omega}_* = \exp\,(\mp\, t/2\tau_c), \qquad \text{...(8)}$$

and

$$\Omega/\Omega_* = 1 \pm (\tau_c/\tau_*)\,[\exp\,(\mp\, t/2\tau_c) - 1], \qquad \text{...(9)}$$

with Ω_*, τ_* beinginitial values.

The last equation shows that Ω/Ω_* would go negative after some $20\tau_c = 6\cdot10^3$ yr for negative *n*, *i.e.* in (roughly) half of all cases. The values in figure 4 cannot, therefore, remain constant for several 10^3 yr.

Two extreme possibilities are worth discussing : First, if the law (8) were valid for time intervals of several 10^3 yr, $\dot{\Omega} \sim \mu_\perp^2/I$ would have to oscillate between its value at birth and much smaller values. This interpretation would explain why spindown ages overestimate true ages, but the timescale for flux decay and buildup—or alignment and counteralignment—would be remarkably short, much shorter than $\tau_\mu (\lesssim 3\cdot10^5\,\text{yr})$ obtained from figure 2.

The opposite and less radical possibility is that the law (8) holds for no longer time intervals Δt than monitored, *i.e.* that the points in figure 4 will change in the immediate future. In this case, $\dot{\Omega} \sim \mu_\perp^2/I$ jitters with amplitude $\Delta t/2\tau_c \approx 10^{-2}$, a jitter which can be fully accounted for by fluctuations in the moment of inertia $I = I(t)$ (instead of by $\mu_\perp$). For this to be true, the weakly coupled shells inside a pulsar would have to possess similar moments of inertia to those inferred from the

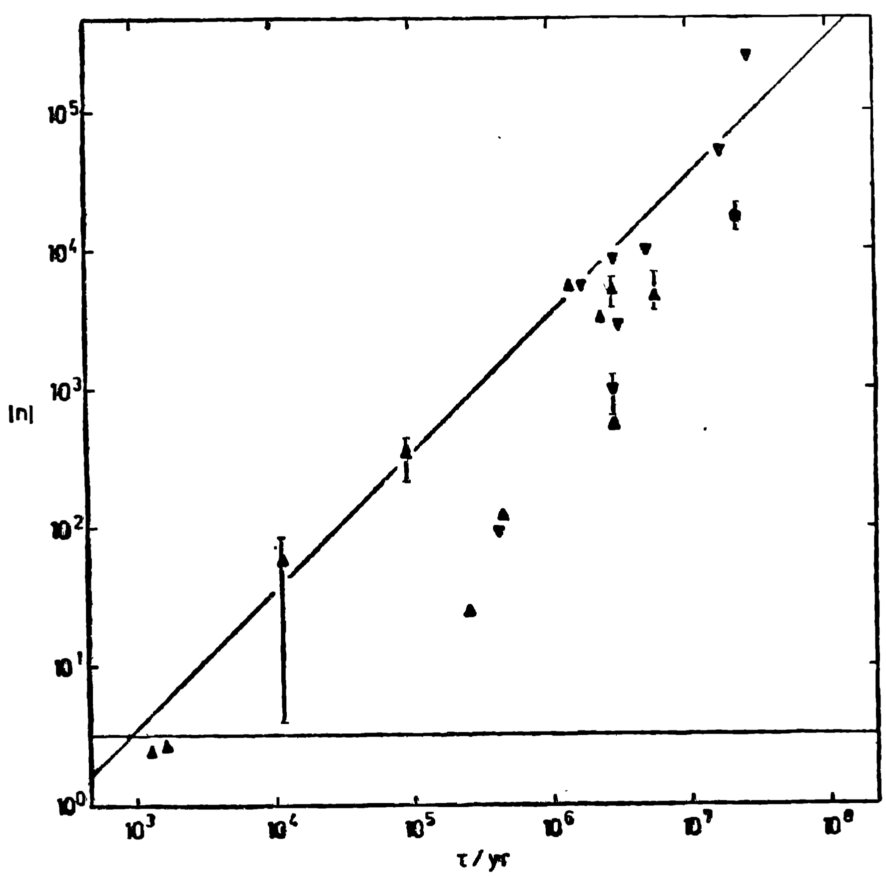

Figure 4. Magnitude of the braking index $n := \Omega\ddot{\Omega}/\dot{\Omega}^2$ plotted *versus* spindown age $\tau := P/2\dot{P}$, whereby triangles 'pointing' upward denote positive values of *n*. $|n|$ is bounded by $|n| \lesssim \tau/3\cdot10^3$ yr instead of being equal to 3 (as predicted for dipole spindown). The data are taken from Gullahorn & Rankin (1982), Downs (1981 : Vela), Demiański & Prószyński (1979), and Manchester *et al.* (1985).

large glitches. The two discussed possibilities were based on the assumption that in the spindown law $\dot{\Omega} = -T/I$, the torque T varies as $\mu_\perp^2$. There is, of course, the logical possibility of yet another coupling to the outside world of which we are so far not aware.

3.4. *The Crab and the Vela pulsar*

For more than 10 years, the 10^3 yr old Crab pulsar was the fastest and youngest pulsar known in our Galaxy. (Table 1 shows that it is no longer the fastest). Its explosion remnant, the Crab nebula, is probably the best studied supernova remnant; see figure 5. Most of its properties are mentioned and/or discussed by Kundt & Krotscheck (1980) and by Davidson & Fesen (1985) where one finds references to earlier literature; see also Kundt (1981b, 1983) for the extended central x-ray source and for its outer edge respectively, and Kennel & Coroniti (1983). Its overall spectrum is shown in figure 6 where dots stand for pulsed emission; recent additions to the spectrum are the infrared 'bump' on the 'elephant's forehead', probably due to ambient dust, and the powerful hard γ's near 10^{16} eV which might be the product of inverse Compton collisions of hard electrons with ultraviolet photons.

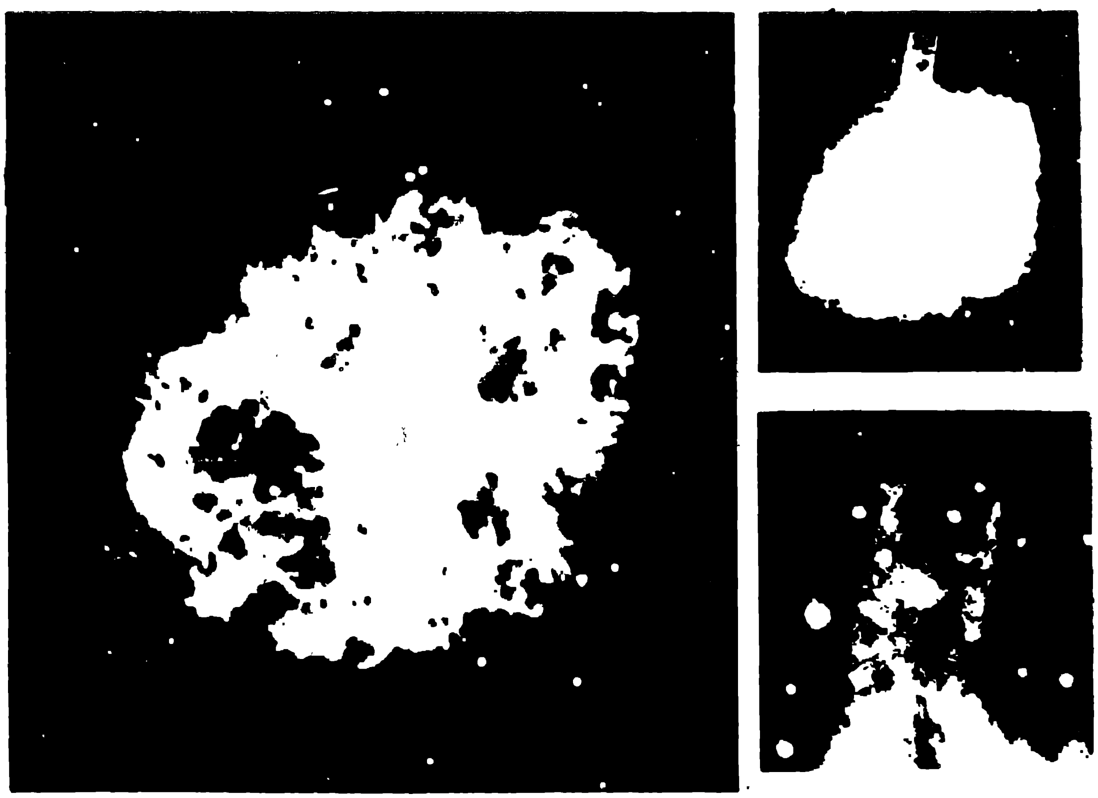

Figure 5. Deep [O III] photograph of the Crab nebula: (*a*) a low-contrast display; (*b*) a very high-contrast display, and (*c*) an enlargement of the northern 'spur'; taken from Gull & Fesen (1982). These emission line photographs suppress the optical synchrotron radiation from the inner part of the nebula. At a distance of $d = 2$ kpc, the long diameter measures 4 pc.

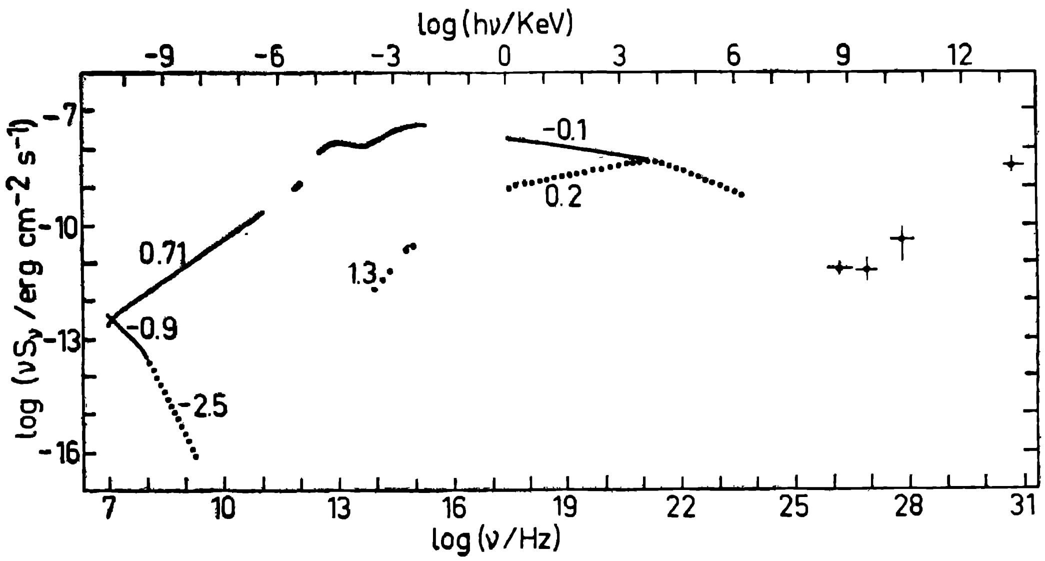

Figure 6. Spectrum of the Crab nebula and its pulsar (dotted), taken from Kundt & Krotscheck (1980) except for the infrared data by Marsden *et al.* (1984) and the (somewhat uncertain) $\gtrsim 10^{16}$ eV point which was communicated at the 18th International Cosmic Ray Conference at Bangalore (1983); see also Graser & Schonfelder (1982: X to γ). Note the wide range of emitted frequencies (photon energies), from $\nu < 10$ MHz to $h\nu > 10^{16}$ eV.

The Crab gives us the unique opportunity to see the interaction of a young pulsar with its nebula. Under the assumption that the bright inner synchrotron sphere, of radius 1 *l*yr (at a distance of 2 kpc), is powered by electrons (and positrons) in transit from the pulsar into the nebula, we can evaluate the injection rate $\dot{N}_{e\pm}$ of relativistic $e\pm$:

$$\dot{N}_{e\pm} = L_{\rm synch}/\langle \gamma \rangle m_e c^2 \approx 4 \cdot 10^{38} \text{ s}^{-1}, \qquad \ldots(10)$$

where $\langle\gamma\rangle = 5 \cdot 10^5 B_{-3}^{-1/2}$ is the average Lorentz factor of the $e\pm$ near the peak of the spectrum, estimated for synchrotron radiation in a magnetic field of $\lesssim 10^{-3}$ G. This injection rate is some 10^4 times the Goldreich-Julian rate $\dot{N}_{GJ}$, which is the rate obtained by extracting from the pulsar magnetosphere its equilibrium electric charge density at the speed of light. $\dot{N}_{GJ}$ is considered an upper bound on the rate of charges to be extracted from the neutron star surface. I conclude that the charges illuminating the nebula are produced in vacuum discharges, probably near the neutron star surface, and are therefore predominantly electrons and positrons, at equal numbers (to within $< 10^{-4}$).

If the Crab is typical, we thus learn that pulsar winds consist predominantly of electrons and positrons. This important conclusion can be corroborated by studies of magnetospheric currents (Cheng & Ruderman 1980), of circular polarization, and of the dynamics of the nebula (Kundt & Krotscheck 1980) whereby protons or heavier ions differ from electrons in their non-degrading inertia (rigidity). Note that independent arguments have been put forward that the extended extragalactic radio sources also are powered by (extremely relativistic) $e\pm$-beams (*e.g.* Kundt 1982a); these arguments have not been immediately taken up by the scientific community (*e.g.* Rees 1981). Their validity (or not) is of great importance for our understanding of the nonthermal sources in the universe.

The Vela pulsar and its nebula Vela X, Y, Z are in many ways similar to the Crab. On the other hand, some distinct differences can serve as counter-examples to suggestive 'extrapolations'. First, the pulsar is not obviously positioned at the centre of its remnant where it should be after no more than 10^4 years, see figure 7. We may learn from this offset that supernova shells are not perfect spheres. Note that the Vela remnant has almost reached its dying age if the 2-component shell model applies; it may turn off in a few 10^3 yr from now.

A second difference between the Vela and Crab pulsars shows up in the pulsed spectrum, and in the phasing of the pulses, see figure 8. The Vela pulses are

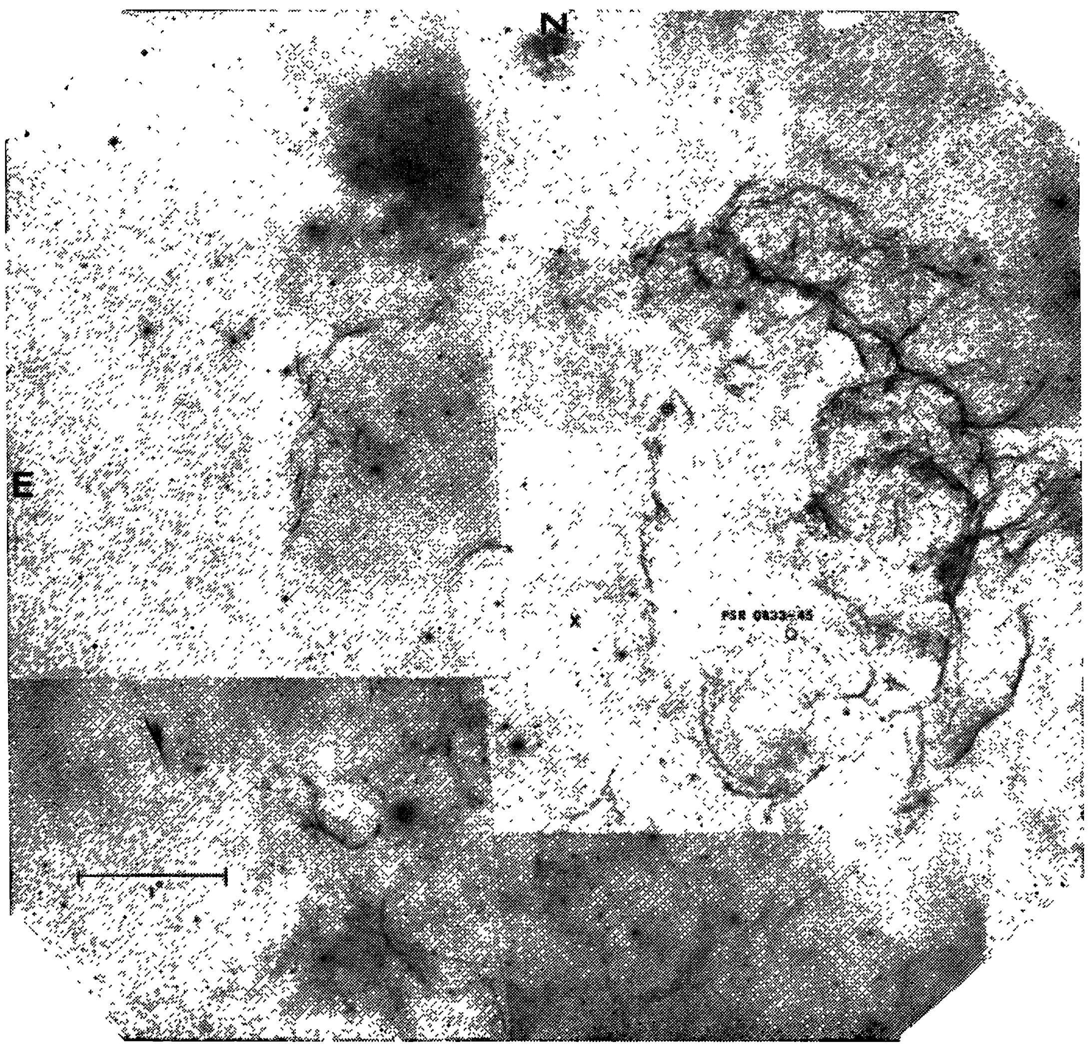

Figure 7. Mosaic of ultraviolet photographs of the Vela supernova remnant taken by Miller (1973), with wavelengths peaking near 3500 Å which are probably due to [O II] emission. The pulsar ought to mark the expansion centre, in which case the complete shell has a diameter of 60 pc (d/0.5 kpc).

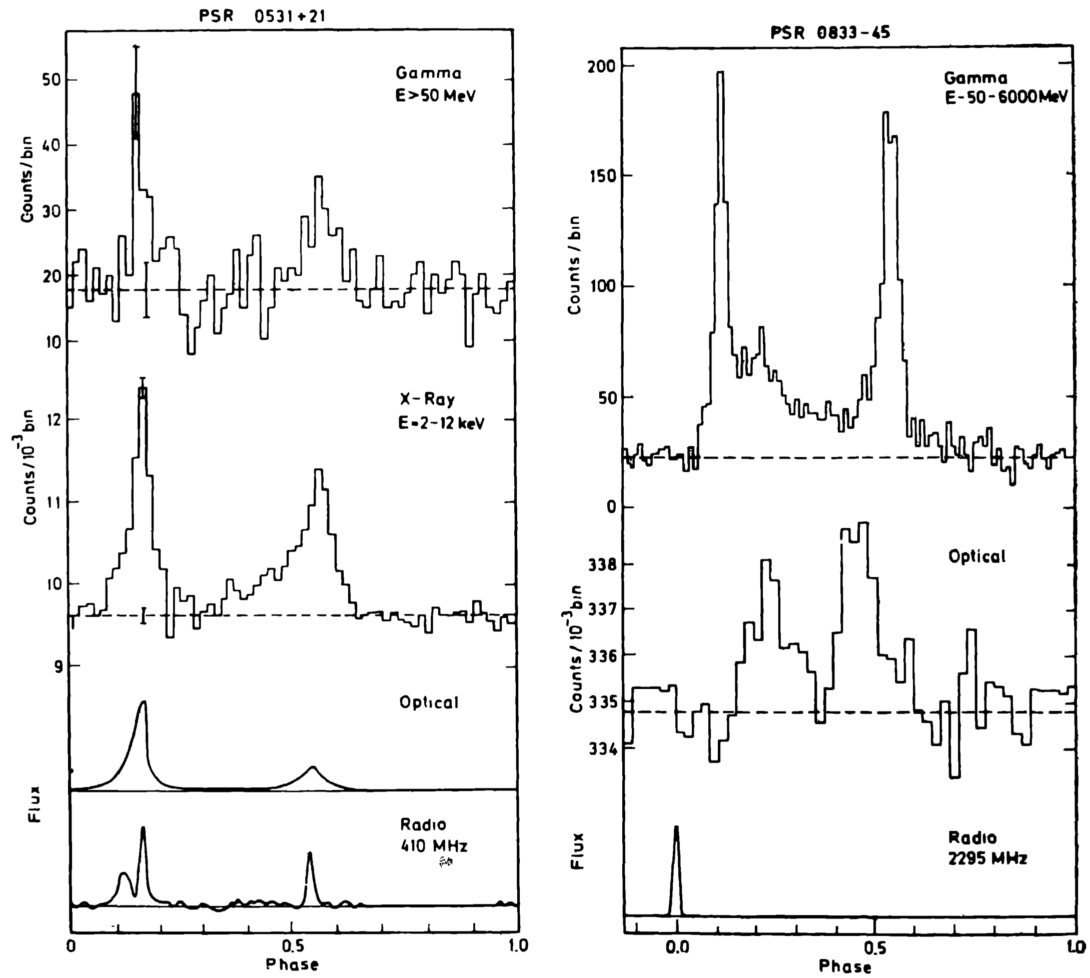

Figure 8. Average pulse shapes of the Crab (0531+21) and Vela (0833–45) pulsars at different frequencies.

stronger at γ-ray energies and yet undetected at x-rays. Only the intensity of the (weak) optical pulses was found as predicted (*cf.* Pacini & Salvati 1983). But the phasing of the two pulses (per period) was unexpected, both absolutely (with respect to the γ-ray pulses) and relatively (pulse and interpulse). If the pulsed radiation is emitted by relativistic $e\pm$ along their instantaneous trajectories, and if these trajectories follow more or less the open magnetic dipole field lines in the sense of an $\mathbf{E} \times \mathbf{B}$ drift, then we learn that the different spectral components must be emitted at different distances from the stellar surface. At most one spectral component could come from near the polar caps, the others must be emitted at distances comparable to, or larger than, the speed-of-light radius $c/\Omega = 3 \cdot 10^8$ cm/Ω.

4. Pulsar models

An extended literature exists on the 'outside story' of a rotating magnetized neutron star, *cf.* Michel (1982), but it is not even clear whether any of the more recently published models by Cheng & Ruderman (1980), Arons & Scharlemann (1979) and

Arons (1981), Michel & Dessler (1983), or others are reliable first-order approximations. It is to my knowledge unknown whether or not there exists a stable steady state (as opposed to an oscillatory ground state) of the rotating pulsar magnetosphere.

4.1. The pulsar magnetosphere

The first and crucial step in pulsar theory was taken by Goldreich & Julian (1969) who pointed out that a rotating magnetized conductor acts like a homopolar inductor. A vanishing electric field in the corotating frame implies an electric voltage ϕ between pole and latitude ϑ (of a homogeneously magnetized sphere of radius R) given by

$$e\phi = eB\Omega R^2 \sin^2 \vartheta/2c \lesssim 10^{18}\ \text{eV}\ B_{13}\Omega_1 R_6^2 \sin^2 \vartheta, \qquad ...(9')$$

where $B_{13} := B/10^{13}\ G$ is the magnetic field in units of 10^{13} gauss, $\Omega_1 := \Omega/10^1\ \text{s}^{-1}$, $R_6 := R/10^6$ cm. This enormous voltage will be almost perfectly shunted by a space charge in the surrounding magnetosphere, because the electric field **E** at the surface has a nonvanishing component parallel to **B** whose force on electrons is some $10^{12} B_{13}\Omega_1$ times larger than gravity.

Force-freeness in the magnetosphere means that wherever the charge density ρ does not vanish, the component of E parallel to **B** must vanish : $0 = \rho \mathbf{E}\cdot\mathbf{B}$, whence $\mathbf{E} = -\boldsymbol{\beta} \times \mathbf{B}$ for some $\boldsymbol{\beta}$. In what follows we assume axial symmetry and stationarity, to simplify the analysis, even though the main conclusions can be generalized to an inclined (oblique) rotator (Mestel 1971). $\boldsymbol{\beta}$ can then be chosen purely toroidal, with $|\boldsymbol{\beta}|$ independent of azimuth ψ, and from $0 = \partial_t \mathbf{B} = -c\nabla \times \mathbf{E}$ one gets Ferraro's isorotation law

$$0 = (\mathbf{B}\cdot\nabla)\, \omega, \qquad ...(10')$$

where ω is the local angular velocity : $\boldsymbol{\beta} =: \boldsymbol{\omega} \times \mathbf{r}/c$. The law (10′) says that along every magnetic field line which is populated by charges, these charges must rigidly corotate with the stellar surface. Corotation can only be violated beyond some vacuum gap if such can form.

Ferraro's law has important consequences. By forming the divergence of $\mathbf{E} = -(\boldsymbol{\omega} \times \mathbf{r}/c) \times \mathbf{B}$, one gets the Goldreich-Julian charge density

$$\rho_{GJ} \approx -\boldsymbol{\Omega}\cdot\mathbf{B}/2\pi c(1 - \beta_{cor}^2), \qquad ...(11)$$

which must be present everywhere in the magnetosphere in order to guarantee corotation. Wherever the charge density deviates from ρ_{GJ} in some sizable domain, electric voltages build up whose strength is of the order of expression (9′). The related number density near the surface corresponds to that in the solar chromosphere : $n_{GJ} = 10^{12}\ \text{cm}^{-3}\ \Omega_1 B_{13}$.

Corotation of charged plasma in a pulsar's magnetosphere is impossible beyond the speed-of-light cylinder, of radius $c/\Omega = 10^{9.5}\ \text{cm}/\Omega_1$, because of the growing inertia of the charges as they approach the speed of light. Centrifugal forces will remove all plasma that gets near there. The lost plasma must be resupplied from inside, in order to avoid deviations from ρ_{GJ}. The magnetosphere thus divides into a closed corotating zone, deep inside the speed-of-light cylinder, and a zone of 'open' magnetic field lines which connect the pulsar to its (distant) surrounding 'nebula', *cf.* figure 9. Whether or not the corotating zone extends out to the speed-of-light cylinder, or perhaps only 10% that far, or whether it exists at all is not known. Numerical approximations show that the corotating zone can only touch the speed-of-light cylinder for a delicate choice of (small) induced surface currents (Wegener 1977), of which it is not known whether they are automatically set up.

In any case, there should be a neighbourhood of field lines from the (magnetic) polar caps which leave the speed-of-light cylinder, along which charges can escape to infinity, launching the 'wind' of the neutron star. These field lines can pass from the positive to the negative space charge sector, determined by $\mathbf{\Omega}.\mathbf{B} \lesssim 0$ respectively (according to equation (11)). Note that the surface $\mathbf{\Omega}.\mathbf{B} = 0$ is a cone for a dipole field $\mathbf{B}$. What currents flow along these open lines ?

At this point, different authors have suggested very different answers. Despite the well elaborated problem of the magnetosphere of the aligned rotator which can be reduced to the so-called 'pulsar equation' for a scalar ('field line') potential (Julian 1973; Scharlemann 1974), it has not been possible to solve the electric current problem from first principles; instead, major difficulties have been pointed out by Okamoto (1975) and others. The different suggested answers range from no wind (for certain field configurations: antipulsars) through favourably-shaped wind channels (slot gap model which contains vacuum sectors) to turbulent halos where the charges escape along some field lines and return to the star along others, after having radiated part of their energy.

4.2. Pulsar wind

Interesting and well motivated as these different solutions are, they do not seem to provide a convincing description of the Crab pulsar which apparently supplies its nebula with $\gtrsim 10^4$ times the Goldreich-Julian density of extremely relativistic electrons, nor do they give a hint as to how Sco X-1 and the active galactic nuclei generate their twin jets. Such strong relativistic winds point at the action of a strong (magnetic) centrifuge, as first analyzed by Gunn & Ostriker (Kulsrud *et al.* 1972; Blandford 1972). A 'strong' electromagnetic wave, *i.e.* a wave which drives charges relativistic so that the Lorentz force exerts a longitudinal push, does not discriminate against one sign of charge, *i.e.* it can drive an electrically neutral wind. Charges reaching the speed-of-light cylinder along open field lines are expected to be seized by the forming wave and swept out almost radially, *cf.* figures 9 and 10.

This Gunn-Ostriker scenario of a strong wave post-accelerating the plasma flow which emerges from the open-field-line region can be supported by a number of further arguments. First of all, the Goldreich-Julian paradox of electric currents flowing from the positive charge sector into the negative one or vice versa can be solved by the presence of a large neutral excess of pair plasma. The condition $\nabla \cdot \mathbf{j} = 0$ of a conserved electric current $\mathbf{j}$ along the open magnetic field lines together with the forcefreeness condition $\rho \approx \rho_{GJ}$ would be violated for a one-component

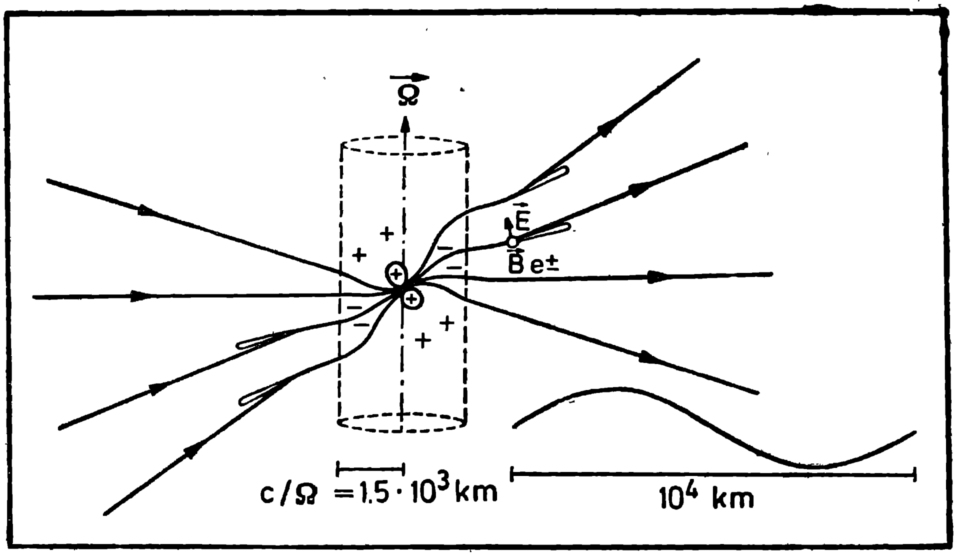

Figure 9. Pulsar at work, in profile (tentative) : An inclined magnet spins around **Ω** and sets up a charged magnetosphere, and near-radial electron-positron flows. The charges radiate predominantly in the transition zone from the corotating magnetosphere (of radius c/Ω) to the distant windzone, whilst they are post-accelerated by the strong (forming) low-frequency wave.

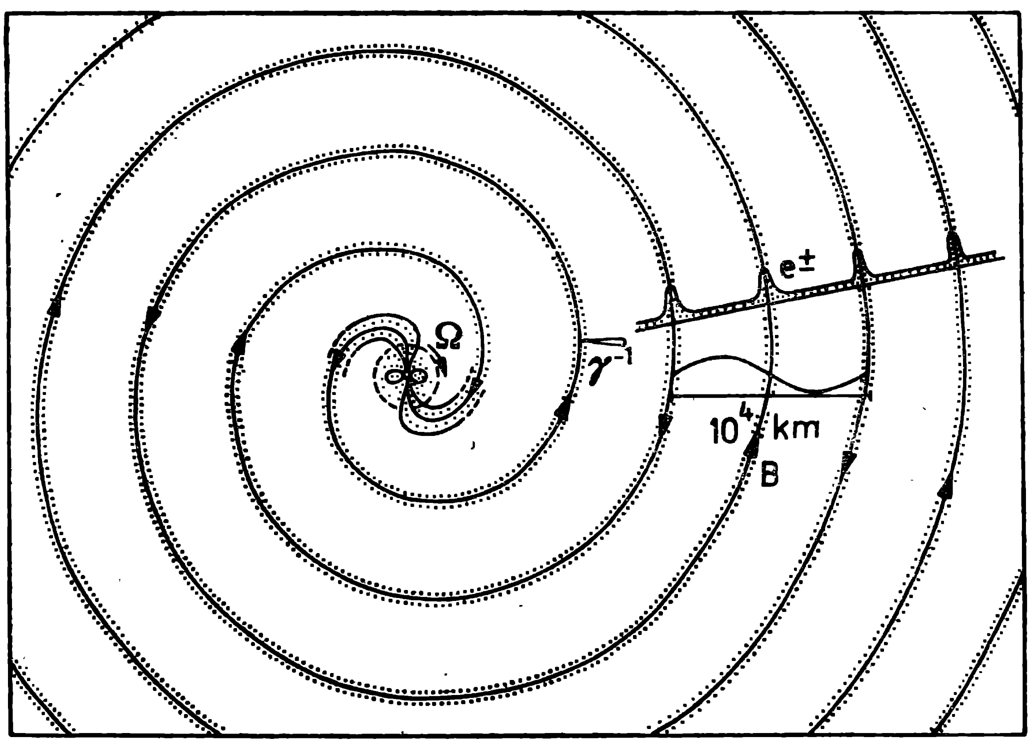

Figure 10. As in figure 9, but seen from the top. The densities of the radially moving electron-positron sheets may correspond to the observed γ-ray light curve.

plasma flowing at (almost) the speed of light, but can be satisfied for a 2-fluid relativistic flow at different Lorentz factors $\gamma_\pm$ if the excess density $n =: \xi\, n_{GJ}$ is large enough to compensate for the smallness in relative velocity :

$$1 \lesssim \xi \mid \beta_+ - \beta_- \mid \lesssim \xi/2\gamma^2_{\min}. \qquad ...(12)$$

In the case of the Crab, a neutral excess of $\xi \gtrsim 10^4$ is consistent with a magnetospheric Lorentz factor of order 10^2 of the slower component.

A second constraint on the wind, to be driven by the magnetic centrifuge, is that its inertia at the speed-of-light cylinder be less than that of the pushing magnetic field :

$$1 \geqslant n\gamma m_e c^2/(B^2/8\pi) \approx 4\xi\gamma/f, \qquad ...(13)$$

where

$$f := \frac{eB}{m_e c\Omega} = 10^9 B_3/\Omega_1 \qquad ...(14)$$

is the 'strength parameter', or nonlinearity parameter of the (forming) wave. Conditions (12) and (13) imply

$$\gamma_{\min} \lesssim f^{1/3}/2 \lesssim 5\cdot 10^2\, (B_3/\Omega_1)^{1/3}, \qquad ...(15)$$

consistent with the earlier constraint. The wave will accelerate the charges until they acquire comparable momentum; according to equation (13), equipartition between wave and particles implies an (asymptotic) bulk Lorentz factor (without radiation losses) :

$$\gamma_{equ} \approx f/4\xi = 10^{4.4}\, B_3/\Omega_1\xi_4, \qquad ...(16)$$

which is comfortably higher than the magnetospheric $\gamma_{\min}$ of the slower component. For the Crab, B_3/Ω_1 is of order 10^3, hence $\gamma_{equ} \approx 10^7$.

A third constraint has been pointed out by Holloway (1975) : If the charges in the pulsar wind perform an $\mathbf{E} \times \mathbf{B}$ drift at essentially the speed of light, their motion is similar to that of a gramophone needle (running backwards) with the grooves representing the magnetic field lines, *cf.* figure 10. One can then compare the voltage difference $\int \mathbf{E}\cdot d\mathbf{x}$ across magnetic field lines in the windzone with that across the same field lines near a polar cap : they must be the same, unless the corotation condition $\mathbf{E} = -\boldsymbol{\beta}_{cor} \times \mathbf{B}$ was violated. In the dipole approximation, the polar caps have an opening angle of $\vartheta_{pc} = \sqrt{\Omega R/c}$. The voltage drop ϕ_{pc} across a polar cap, or rather its corresponding Lorentz factor γ_{pc}, then follows from equation (9′) :

$$\gamma_{pc} = e\phi_{pc}/m_e c^2 = (eBR/m_e c^2)\,(\Omega R/c)^2 = 10^{8.3} B_{13}\Omega_1^2. \qquad ...(17)$$

Holloway's constraint demands that the total voltage drop along an open field line be small compared to ϕ_{pc}.

This constraint becomes somewhat more severe when we remember Michel's relativistic generalization of Child's law for a space-charge-limited flow. From the conservation laws of electric charge ($\nabla \cdot \mathbf{j} = 0$) and energy ($e\delta\phi = \delta\gamma m_e c^2$) together with Maxwell's equation $\nabla . \mathbf{E} = 4\pi\rho$ one finds for a stationary, 1-d, 1-fluid fall through the potentiol $\phi =: \gamma m_e c^2/e$ the asymptotic Lorentz factor γ_∞,

$$\gamma_\infty \approx (8\gamma)^{1/2}, \qquad \text{...(18)}$$

cf. Sutherland (1979). When applied to polar cap acceleration, this formula leads to the constraint

$$\gamma_\infty \lessdot 4 \cdot 10^4 B_{13}^{1/2} \Omega_1, \qquad \text{...(19)}$$

corresponding to energies $\gamma_\infty m_e c^2 \lessdot 2 \cdot 10^{10}$ eV, (still) large enough to guarantee ample production of $e^\pm$ pairs. Note that photoproduction of $e^\pm$ in a transverse field $B_\perp$ needs

$$h\nu/m_e c^2 \gtrsim m_e^2 c^3/B_\perp e\hbar = 1/B_{\perp 13 \cdot 6}, \qquad \text{...(20)}$$

and pair production by photon-photon collision (of energies E_j) needs

$$E_1 E_2 \geqslant (0.511 \text{ MeV})^2. \qquad \text{...(21)}$$

Both conditions are easily met near the polar caps of a young neutron star, and, more generally, near the polar caps of a sufficiently fast rotating strong magnet.

How large is the electric current density j in a pulsar wind? Many authors worry about return currents, to ensure charge neutrality of the neutron star. Cannot j vanish, *i.e.* be small compared with the Goldreich-Julian value $\rho_{GJ}c$? As long as there is a large excess pair plasma, $\xi \approx 10^4$, this plasma will endeavour to shunt large excess voltages. Such voltages would in particular build up near the magnetic polar caps if j were of order $\rho_{GJ}c$, because current loops cannot close radially across the caps, due to the perpendicular magnetic field; instead, we would have toroidal Hall currents. Even if radial (Pedersen) currents could flow across the polar caps, they would tend to brake the neutron star's spin with a very short lever arm. The polar caps would thereby be strongly sheared and heated and become bright x-ray sources, brighter than observed Finally, current densities comparable to $\rho_{GJ}c$ would twist the magnetospheric field, giving rise to torques whose reaction needs a higher inertia than provided by the (dynamically insignificant) pair plasma. For all these reasons, I expect an almost vanishing current density j in a pulsar wind.

4.3. Strong wave

If only insignificant currents flow in a pulsar windzone, the spindown torque must be exerted through the emission of strong magnetic dipole waves. For the vacuum case, the electromagnetic fields have been calculated in closed form by Deutsch (1955). Their modification by the outflowing magnetospheric plasma is unknown.

In an attempt to study a strong plasma wave, Akhiezer & Polovin (1956) have cast the relativistic equations of motion (without radiation reaction) for the particles and fields of a plane wave into Lagrangean form. More generally, an arbitrary 1-d plasma wave satisfying $F_{[ab}k_{c]} = 0$ can be treated in this way, where $F_{ab}(\varphi)$

is the field strength tensor, $k^a = (\omega/c, \mathbf{k})$ is the 4-d wave vector, $\varphi := k_a x^a$ is the phase of the wave, $j^a = c(\rho_+ u_+^a + \rho_- u_-^a)$ its 4-current, $\rho_\pm$ = charge densities of the (2-component) plasma and $u_\pm^a u_{\pm a} = 1$. When the field strength tensor is expressed through a 4-potential $A_b(\varphi)$ as $F_{ab} = 2A_{[b,a]}$ and A_b satisfies the (transverse) radiation gauge $A_b k^b = 0$, all quantities can be formulated as functions of the 'transverse' 3-potential **A**. **A** satisfies the equations of motion of a Newtonian mass point in an axially symmetric potential $V(\mathbf{A})$. Its motion can be approximated analytically for large and small $|\mathbf{A}|$, and solved numerically throughout (*cf.* Kundt 1985c).

In this program, one has to distinguish between waves with a timelike or spacelike 4-vector k_a, called 'superluminal' or 'subluminal' respectively. For superluminal waves there is a preferred frame in which the wave is spatially homogeneous and the magnetic field vanishes. In this frame, **A** is an arbitrary spacelike 3-vector, and $V(\mathbf{A})$ is a stable potential which acquires spherical symmetry for large $|\mathbf{A}|$. In this formulation it becomes obvious that circularly polarized large-amplitude waves are stable whereas transverse linearly polarized ones are not (*cf.* Che & Kegel 1980).

For subluminal waves, on the other hand, **A** is formally a vector in a 3-space of hyperbolic metric, $V(\mathbf{A})$ has saddle shape, and moreover there are forbidden regions in phase space. For this reason, subluminal waves have acquired little popularity even though they correspond intuitively to the joint motion of relativistic particles and electromagnetic fields at marginally relativistic speeds (Goldreich & Julian 1970, Kundt & Krotscheck 1980). Note that no strictly periodic behaviour is expected when the wave transfers (part of) its 4-momentum to the charges, and that no global solutions are required because a 1-d wave can at best be a local approximation to the pulsar wavezone.

Apart from this somewhat unrealistic nature of plane waves one can estimate their damping time through the emission of hard γ-rays. As long as radiation losses

$$dE/dt = (2e^2/3c)\, \dot{u}^a \dot{u}_a = (2e^4/3m^2c^3)\, F_{ab}u^b F^{ac} u_c$$

are small, such losses can be evaluated in closed form. The corresponding particle-energy e^{-1}-folding time t_{rad} (in the lab frame) reads, in units of the wave period $2\pi/\omega$,

$$\omega t_{rad} = \frac{3m_e c^3}{2e^2 \omega \tilde{f}^2 \gamma} = \frac{1}{\omega_1 \tilde{f}_9^2 \gamma_4}, \qquad \ldots(22)$$

where (Kundt 1985c)

$$\tilde{f}^2 := -(e/m_e c\omega\gamma)^2 F_{ab}u^b F^{ac}u_c = (e/m_e c\omega)^2 [(\mathbf{E} + \boldsymbol{\beta} \times \mathbf{B})^2 - (\mathbf{E}\cdot\boldsymbol{\beta})^2] \qquad \ldots(23)$$

is of the order of the squared strength parameter f^2 introduced in equation (14), with $0 \leqslant \tilde{f} \lesssim f$. For waves with $E > B$, the inequality

$$|F_{ab}u^b F^{ac}u_c| \geqslant |F_{ab}F^{ab}|/2 = E^2 - B^2 = E^2(1 - c^2k^2/\omega^2) \qquad \ldots(24)$$

implies $\tilde{f} \geqslant f(1 - \beta_{wave}^{-2})^{1/2}$, so that damping is rapid unless $\beta_{wave} \approx 1$ ($k_a \approx$ light-like). $\tilde{f}$ can only vanish in the subluminal case, namely for $F_{ab}u^b = 0$, *i.e.* for

$\mathbf{E} \times \mathbf{B}$-drifting charges (MHD approach). Note that magnetic dipole radiation has $E < B$.

When equation (22) is compared with equations (14) and (16) it tells us that radiation losses are in general fast compared with the wave period. Thus the radiation term must not be omitted from the equations of motion except for $\tilde{f} \approx 0$ in the subluminal case; *cf.* a similar conclusion by Asséo *et al.* (1978). Moreover, there is the Weibel instability of mutual attraction of parallel currents which causes even higher radiation losses (Asséo *et al.* 1980). In contrast, the Crab spectrum shows only little very hard radiation (a few per cent of the total power), which may even be due to inverse Compton losses (figure 6). I conclude that a strong wave with a high plasma load ($\xi \approx 10^4$) is not well approximated by a superluminal wave.

4.4. *Pulsar radiation*

This section must be even scantier than the ones before because no deep understanding appears yet to have been achieved. Many authors have envisaged source locations deep inside the magnetosphere which have difficulties with the different phasing at different frequencies (figure 8) and also with the apparently large beaming factors ($f \gtrsim 0.5$: Narayan & Vivekanand 1983) and occasional large pulse windows ($< 1\%$ to $> 90\%$). If the charges move relativistically from near the neutron star surface to the outer shock of the pulsar windzone, then radiation emitted anywhere along this path will be observed as pulsed emission, as long as there is not a large dispersion of particle directions so that we average over emission spheres of non-negligible spherical extent. The mere fact that the radiation is pulsed does not tell us from where it is coming. Rather, the phase offset $\Delta\varphi$ of interpulses (of order $\Delta\varphi/\varphi \approx 0.1$) may measure the distance of the emission site.

With this proviso in mind, the post-acceleration zone between the corotating magnetosphere and outer windzone is a plausible site for emission because it involves the largest accelerations. Note that radio emission from near the speed-of-light cylinder is not necessarily excluded by scintillation experiments which limit the transverse source extent to $\lesssim 10^3$ km (corresponding to $\lesssim 1\%$ of the speed-of-light radius : Cordes *et al.* 1983) as we only sample over those (small) subareas of the emission region where particle trajectories point towards the earth within an angle of γ^{-1}. Another theoretical argument against near-zone emission is the enormous optical depth τ of the outer magnetosphere for induced Compton scattering (towards lower frequencies) :

$$\tau \gtrsim \sigma_T \int_{r_{em}}^{\infty} n ds \, (kT_b(\nu_{min})/m_e c^2) \, (\nu_{min}/\nu)^{1/2} \, \gamma^{-4} \approx (\nu_{min}/\nu)^{1/2} \, \gamma_4^{-4} \qquad \ldots(25)$$

which makes the escaping wind opaque to radio waves as long as it contains electrons with Lorentz factors $\gamma < 10^4$ (Wilson & Rees 1978). In their discussion of curvature radiation by pulsars, Ochelkov & Usov (1980) seem to bypass this difficulty by assuming large enough Lorentz factors.

For the sake of argument let us fit the radiation from the Crab pulsar and others to synchrotron radiation emitted near the speed-of-light cylinder, whereby we leave the transverse excursions (pitch angles α) of the electrons as an open parameter in the form of an unknown transverse field strength $B_\perp =: B \sin \alpha$. Synchrotron radiation peaks at a frequency

$$\nu_s \approx \gamma^2 e B_\perp/\pi m_e c \approx 10^{14}\ \text{Hz}\ \gamma_4^2\,(B_\perp/0.2\ G). \qquad ...(26)$$

The incoherent synchrotron power L_{incoh} emitted at a distance r from the star can then be written in the form

$$L_{\text{incoh}} \approx \gamma^2 n r^2 \Delta r e^4 \widetilde{B}^2/m_e^2\, c^3 \qquad ...(27)$$

$$\approx 3\cdot 10^{33}\ \text{erg s}^{-1}\ (\widetilde{B}^2/B_\perp\cdot 10^6\text{G})\ (\nu_s/10^{14}\text{Hz})\ \mu_{31}\xi_4\Omega_2,$$

where $\widetilde{B}^2 := (\mathbf{E} + \boldsymbol{\beta} \times \mathbf{B})^2$, $\boldsymbol{\beta} := v/c$, $\Delta r =$ thickness of the radiating layer $\lesssim r$, and $\mu := Br^3 =$ dipole moment. It is of the correct size of the infrared pulses from the Crab for $\widetilde{B}^2/B_\perp \lesssim 10^6$ G, and scales as

$$L_{\text{incoh}} \sim \gamma^2 n r^3 B^3 \sim \gamma \mu^4 r^{-9} \sim \gamma \mu^4 \Omega^9 \qquad ..(28)$$

with evolving angular frequency Ω if we assume that the emission distance r scales as the speed-of-light distance c/Ω, $B \sim r^{-3} \sim \Omega^3$, and that inequality (13) scales like an equality : $\gamma n \sim B^3$. Essentially this scaling law was first advanced by F. Pacini. If moreover ξ is independent of Ω, we get $\gamma \sim \Omega^2$, $L_{\text{incoh}} \sim \mu^4\Omega^{11}$, and $\nu_s \sim \Omega^7 \times \sin\alpha$. These proportionalities are consistent with our knowledge about the optical and infrared pulses from (fast) pulsars, though some modification (by beaming?) has to be invoked for the 50 ms LMC pulsar (Middleditch & Pennypacker 1985).

The radio pulses (of frequency ν) have such a high brightness that they are thought to be emitted coherently, by bunches of N particles oscillating in phase. Their linear extent cannot be much larger than $c/2\pi\nu$ so that we have

$$N \approx n(c/2\pi\nu)^3. \qquad ...(29)$$

Alternatively, N can be expressed through the brightness temperature T_b or the blackbody spectral intensity (at the photosphere) $B_\nu = kT_b c^2/2\pi\nu^2$ in the form

$$N \approx kT_b/\gamma m_e c^2 = (S_\nu/2\pi\nu^2\gamma m_e)\,(d/r)^2, \qquad ...(30)$$

where d is our distance from the source, and $S_\nu = B_\nu(r/d)^2$. By equating both expressions for N we get the condition

$$\gamma n r^2 \approx 4\pi^2\nu S_\nu d^2/m_e c^3 = 10^{27}\ \text{cm}^{-1}\ L_{31} \qquad ...(31)$$

for the typical radio power $L_R = 4\pi d^2\nu S_\nu = 10^{31}$ erg s^{-1} in subpulses near $\nu =$ 1 GHz; (observed instantaneous values of L_R range between 10^{27} and 10^{32} erg s^{-1}). Insertion into the first expression for the coherence factor N yields

$$N = 10^8/r_9^2\,\nu_9^3\,\gamma_3, \qquad ...(32)$$

a reasonable number. Finally, on insertion of equation (29) into equations (27) and (26) one gets for the coherent radio power $L_{\text{coh}} = NL_{\text{incoh}}$

$$L_{\text{coh}} \approx n^2 r^2 \Delta r c e^3 \widetilde{B}^2/2(2\pi\nu)^2\, m_e B_\perp$$

$$\approx 10^{31}\ \text{erg s}^{-1}\ (B\widetilde{B}^2/B_\perp 10^4 G^2)\ \Omega_1^2\,\xi_4^2\,\mu_{31}/\nu_9^2, \qquad ...(33)$$

which involves much lower values of $\tilde{\mathbf{B}}$ than the expression for L_{incoh} when fitted to typical observed luminosities (which can, however, be temporarily exceeded by large factors). Note that L_{coh} formally decreases with distance r from the pulsar, *i.e.* favours emission sites near the star if $\tilde{\mathbf{B}} = \mathbf{E} + \boldsymbol{\beta} \times \mathbf{B}$ does not vanish there, yet is by far large enough to allow emission sites in the vicinity of the speed-of-light cylinder.

These phenomenological radiation expressions fall very short of explaining the wealth of observed detail. For instance, radio pulses are composed of subpulses, of typical duration ms, superimposed by needle-sharp micropulses, of duration down to and less than 0.5 μs, with $\Delta P/P \leqslant 5 \cdot 10^{-7}$ (Hankins & Boriakoff 1978). Not all pulsars show such extreme microstructure (containing a significant fraction of the total radio power), at least not at frequencies above 1 GHz, but the fact that some of them do show how bright pulsar radiation can be. Next there is the phenomenon of drifting subpulses, and of pulse nulling, different from pulsar to pulsar. Fourth, the intensity distribution of the Crab radio pulses below 500 MHz has a long power-law tail towards giant pulses, with upto $\gtrsim 10^3$ times the typical strength, being distributed according to $dN \sim I^{-3.5}\, dI$; see also Ritchings (1976) for other pulsars. Fifth, pulse windows tend to open up towards 180° when the dynamic range is increased beyond $10^2:1$.

A further remarkable structure of the radio emission of many pulsars is that several of their properties change at the same characteristic frequency ν_b, of order several GHz: The power-law spectra often show a break towards softer spectra at higher frequencies; there is a break in the power-law distribution of the separation and width of subpulses, in the opposite sense, and similarly in the modulation index, which measures the energy scatter in successive pulses. Moreover, microstructure disappears at high frequencies. Remarkably, all of these breaks occur at the same frequency, pointing at a qualitative change in the coherence mechanism (Bartel *et al.* 1980).

Another outstanding characteristic of pulsar radiation is the high degree of linear and circular polarization, often reaching 100% in subpulses, and the systematic change in the position angle of linear polarization which can be used to estimate both the shape (elongated in latitude and curved) and orientation of the pulsar's antenna lobe (Narayan & Vivekanand 1982, 1983). Finally, there is the systematic distribution of pulsar profiles over the P, $\dot{P}$-plane, or $B_\perp,\tau$-plane, which (also) indicates a stretching of the hollowcone beam pattern in latitude direction, and the presence of a non-drifting, circularly polarized soft core emission inside the harder, linearly polarized hollow-cone (Rankin 1983). An observation of many pulsars allows a statistical reconstruction of (the latitude structure of) their antenna pattern.

5. Binary neutron stars

If neutron stars form from massive stars, and if massive stars are born as double (or multiple) star systems, as discussed in section 2.3, then we expect roughly as many young binary neutron stars in the Galaxy as single ones. In this conclusion, it is assumed that the first supernova explosion in a binary is unlikely to disrupt the system. This assumption is based on the expression for the (numerical) eccentricity ϵ of a binary

$$\epsilon = (M_1^- - M_1)/(M_1 + M_2) \qquad ...(34)$$

after an instantaneous, isotropic, and interaction-free removal of some mass $M_1^- - M_1$ from the first star, whereby M_1 and M_2 are the final masses, and the initial orbit has been assumed circular. A system stays bound as long as $\epsilon < 1$, *i.e.* as long as less mass is lost during the explosion than remains. In a massive star system, this condition is expected to be satisfied for the first explosion but not for the second, so that the first-born neutron star tends to be binary whereas the last-born one tends to be single, hence a pulsar.

Indeed, we are aware of (massive) runaway stars which may well have a neutron star companion; they may comprise as many as 50% of all massive stars (Stone 1979, and section 2.3). We also observe some 150 bright Galactic x-ray sources ($L_x \gtrsim 10^{35}$ erg s^{-1} in the energy range 1–10 KeV) which are held to be powered by mass accretion onto a neutron star; apparently the bright x-ray stage of massive binaries is short-lived. It needs a high masstransfer rate and may involve slowly spinning neutron stars, with periods in the second to minute range, as observed for the pulsing sources. If first-born neutron stars have similar initial spin rates ($\lesssim 10^{-1}$ s) to their younger brothers, the pulsars, then their spindown must be difficult to observe even though it involves the liberation of an enormous rotational energy. Are we blinded by the bright companion?

In addition to the bright x-ray sources whose intensity tends to flicker on time-scales between several ms and yrs, and is sometimes pulsed ($\approx 20\%$), there are the less bright Galactic x-ray sources—which are often, but not always, white-dwarf systems—the long-lived x-ray bursters habitating the Galactic bulge which involve low-mass companions, the γ-ray bursters which may be old, nearby neutron stars (Lipunov *et al.* 1982), and finally some four binary radio pulsars which are probably a rare subclass of the pulsars ($\gtrsim 1\%$).

All these objects are interesting sources on their own, but their main importance for our understanding of neutron stars is that they give us a chance to estimate neutron star masses, radii, moments of inertia and magnetic dipole moments. We shall discuss them more or less in evolutionary order.

5.1. *The dormant stage*

We are not strictly aware of fast spinning neutron stars with massive companions, though the runaway stars may be such systems. Further candidates are the jet sources Sco X-1 and SS 433 whose accreting neutron stars may have a spin rate comparable to that of the Vela pulsar (on energetic grounds). Yet the new-born Crab had a rotational energy $E_{rot} = I\Omega^2/2 = 10^{49.5}$ erg which can supply the Eddington luminosity of some 10^{38} erg s^{-1} for 10^4 yr, during which time we should be able to see it throughout the Galaxy if that luminosity were radiated in monitored frequency bands. An even some 10^3 times larger rotation energy is stored in the 1.56 ms pulsar because the gravitational binding energy is reduced in magnitude for large Ω (Cowsik *et al.* 1983).

On the other hand, the pulsing x-ray sources show extremely low spin rates. Unless they were born slow, their rotational energy at birth must have been transferred somewhere else. We do not observe the equivalent radiation; and even if we did, there is the problem of the missing angular momentum : Transport of angular momentum at relativistic speeds would need an average lever arm of $c/\Omega = 10^{8.5}$ cm$/\Omega_2$.

Most likely, the angular momentum has been transferred to the wind matter of the companion star, in which case a braking mass ΔM smaller than $10^{-4.8}\ M_\odot$ gets accelerated to relativistic velocities (by conservation of $10^{49.5}$ erg). A typical stellar mass loss rate of some $10^{-9}\ M_\odot\ \mathrm{yr}^{-1}$ during 10^5 yr, of which some 10^{-2} are attracted by the neutron star, would correspond to a total braking mass ΔM of order $10^{-6}\ M_\odot$ with a resultant average Lorentz factor above 10.

In other words, the fast spinning neutron star may act as a booster of cosmic rays, like a relativistic grindstone (Holloway *et al.* 1978; Kundt 1982b). Figure 11 is a drawing of the expected geometry during this stage, which will end as soon as an accretion disc forms and confines the corotating magnetosphere deep inside its speed-of-light cylinder.

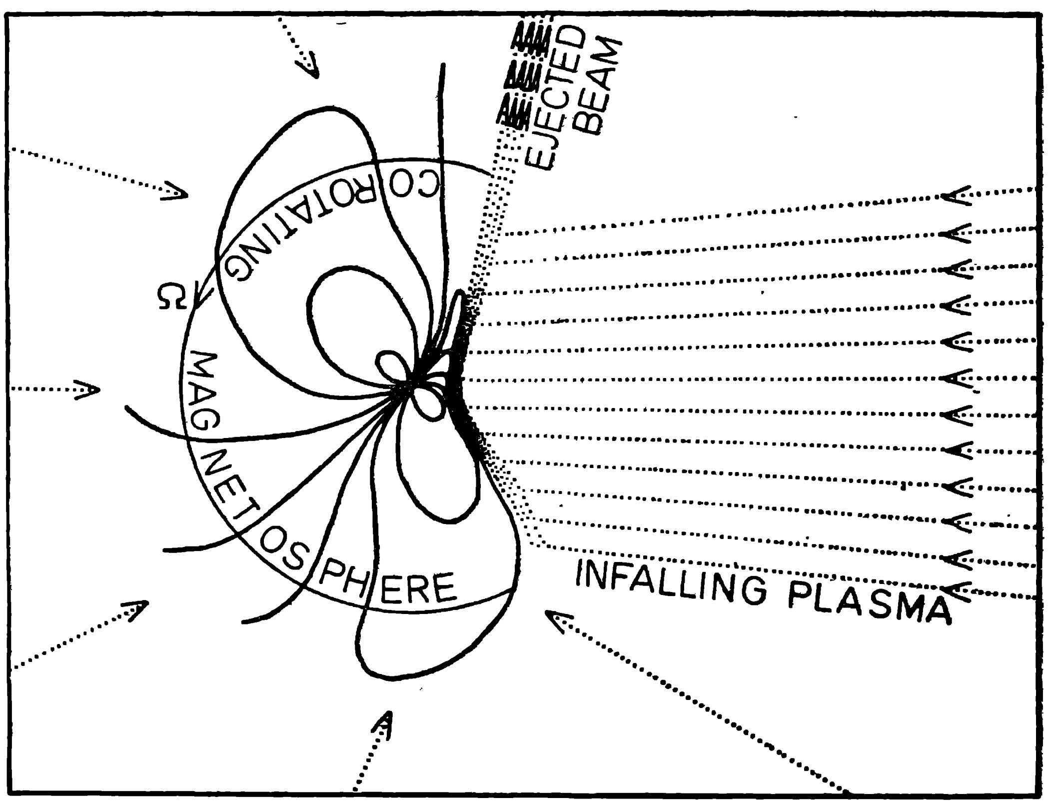

Figure 11. Tentative geometry of the grindstone stage, during which an infalling wind-plasma from the companion star confines the corotating magnetosphere and is ejected in the form of a cosmic-ray beam, thereby reducing the neutron star's spin (*cf.* Kundt 1982*b*).

5.2. *The pulsing binary x-ray sources*

The pulsing binary x-ray sources—sometimes called x-ray pulsars for short—form a minority among the bright Galactic x-ray sources : only 24 are currently known; see table 2. All the other bright x-ray sources are unpulsed, with $\Delta I/I < 1\%$. For a recent review see Rappaport & Joss (1983); the properties of the optical counterparts are presented by Bradt & Mc Clintock (1983). In particular, mass determinations of six x ray pulsars plus the binary (radio) pulsar from their Keplerian motion give masses which are all consistent with $M = (1.4 \pm 0.2)\ M_\odot$. Four further neutron star candidates (Sco X-1, Cyg X-2, 3 U 1700–37, and the companion of the binary pulsar) are likewise consistent with this mass range.

It is much harder to derive the neutron star's magnetic dipole moment from the known data (Wasserman & Shapiro 1983). Note, however, that many of the pulsing neutron stars tend to be much older than the critical $3 \cdot 10^5$ yr beyond which the pulsar distribution shows a decline of $\mu_\perp$ (section 3.1), because their companion has first to leave the main sequence and expand (van den Heuvel 1983). It is therefore of considerable importance to do the analysis.

A first item to be clarified is whether or not the x-ray pulsars are wind-fed or disc-fed, that is whether or not the accreting material has enough angular

Table 2. Table of all the known (26) pulsing binary x-ray sources, in order of increasing spin period P. P_{orbit} = orbital period; L = x-ray luminosity; $\dot{P}$ = time derivative of spin period; $< \dot{P} >$ = long time average ($\gtrsim$ several months) of $\dot{P}$; and $Q := L\Omega^{-7/3}/10^{38.5}$ erg $s^{4/3} \approx \xi/\mu^2_{31}$ measures the coupling efficiency between the star and disc. In the four cases where two values for $P/< -\dot{P} >$ are given, the value in parentheses is characteristic of the longest observed epoch. $d_{(1)}$ stands for distance/10 kpc. The data are taken from Bradt & Mc Clintock (1983), Bignami *et al.* (1984), Stella *et al.* (1985), and Hayakawa (1985).

Source	P/s	P_{orbit}/d	$L/10^{37}$ erg s^{-1}	$P/<-\dot{P}>$ yr	Q
A 0538−66	0.069	16.66	80	−4 ?	0.7E−4
SMC X−1	0.71	3.892	60	1.4E3	1.2E−2
Her X−1	1.24	1.700	1	3.4E5	0.8E−3
4U 0115+63	3.61	24.31	1	3 E4 ?	0.9E−2
V 0332+53	4.38	34.25	$0.07d_{(0.2)}^2$	> E4	E−3
Cen X−3	4.84	2.087	5	3.4E3	0.9E−1
1E 2259+586	6.98	0.03	0.02		0.9E−3
4U 1627−67	7.68	0.0288	$08d_{(1)}^2$	5 E3	4.3E−2
2S 1553−54	9.26	30.7 ?			
LMC X−4	13.5	1.408	35	4 E3	0.7E 1
2S 1417−62	17.6	>15	⩽4		⩽1.5E 0
0A0 1653−40	38.2		>0.04	>2 E2 ?	>0.9E−1
1E 0630 + 178	60		$0.00001d_{(\bullet)}^2$	< E3	0.7E−4
4U 1700-37	67.4	3.4	0.3		2.4E 0
A 0535+26	104	111	⩽2	E3	⩽4.7E 1
GX 1+4	122	>15	4	4.7E1	1.4E 2
4U 1230−61	191				
GX 304−1	272	135 ?	0.2		4.4E 1
Vela X−1	283	8.965	0.15	⩾3 E3(E5)	3.6E 1
4U 1145−61	292	187 ?	0.03	⩾ E3(∞)	0.8E 1
1E 1145.1−6141	297	>12	0.3	⩾3 E2(∞)	0.8E 2
A 1118−61	405		0.5		2.8E 2
4U 1907+09	438	8.38	4	> E2	2.5E 3
4U 1538−52	529	3.730	0.4	⩾5 E2 ?	4.1E 2
GX 301−3	696	41.4	1	⩾ E2(3E3)	2.0E 3
X Per	835	580 ?	0.001	⩾1.4E3	3.0E 3

momentum to be temporarily stored in an accretion disc. There are indications of the presence of an accretion disc in at least five of the sources (van Paradijs 1983). The evidence is based on details of the light curve, on weak emission lines (Fe, He II, N III), on the optical-to-x-ray continuum, and on the presence of periods other than spin and orbital. For the remaining sources (which tend to have the longer spin periods P), no direct evidence is available.

But their spinup curves $P(t)$ look almost indistinguishable when drawn to proper scale. No case of secular spindown is known after $\gtrsim 5$ years of monitoring: 9 of the 26 pulsars clearly spin up, the others are marginal. The excursions of the spin periods around their momentary average values look highly non-sinusoidal, like garlands with periodicities between days and years, and they are often—but not always—correlated with the x-ray intensity (*e.g.* Howe *et al.* 1983). Spinup is often interrupted by spindown, with $\Delta P/P$ ranging between $2 \cdot 10^{-2}$ (GX 1 + 4) and $3 \cdot 10^{-6}$ (Her X-1), such that intermittent spinup- and spindown-rates $|\dot{P}|$ exceed the average $\langle -\dot{P} \rangle$ by factors of 3 to $\gtrsim 30$ in most of the sources. The largest excursions satisfy the empirical law

$$|\Delta P| \approx \langle -\dot{P} \rangle \text{ yr}. \qquad \ldots(35)$$

Table 2 shows that the periods, spinup times and coupling efficiencies Q are distributed quasi-continuously over the 26 objects, perhaps with a period gap between the fastest and the second fastest. Consequently, there is no clear borderline between wind-fed and disc-fed sources. Moreover, the local asymmetry between spinup and spindown and the detailed correlations between $\dot{P}$ and the x-ray intensity pose problems to wind-feeding, and even more so to two different mechanisms for which all these characteristics are similar; see also Burnard *et al.* (1983) and Lipunov (1982).

I therefore restrict the following considerations to disc accretion, the only scenario for which the neutron star's dipole moment can (so far) be assessed. The instantaneous geometry may then look similar to figure 12, in which the disc has been approximated infinitely thin and diamagnetic (Kundt & Robnik 1980). In reality, field lines will somewhat cut into the inner edge of the disc and be dragged along in toroidal direction by the Keplerian disc material. The disc material spirals inward through the disc and probably crosses the inner edge in clumps, via Rayleigh-Taylor instabilities, until it reaches field lines that lead directly 'down' to the polar caps. Dense clumps may alternatively survive all the way down to the surface, in the form of metallic blades (Kundt *et al.* 1985). Part of the material may also be 'scratched off' the faces of the disc and be flung out again, in the form of a wind.

Where exactly is the inner edge of the disc? When the disc first forms and spreads inward (and outward) from some accretion torus, it cannot discharge as long as its inner edge is outside the corotation radius r_{cor} (where gravity is balanced by centrifugal forces) because charges entering the corotating magnetosphere are flung out again. Accretion can only start when the inner radius a has shrunk to r_{cor}, where

$$r_{cor} = (GM/\Omega^2)^{1/3} = 6 \cdot 10^8 \text{ cm } \Omega_0^{-2/3} \qquad \ldots(36)$$

holds for $M = 1.4\, M_\odot$, $\Omega_0 := \Omega/10^0 \text{ s}^{-1}$. Suppose accretion through the inner edge is slow so that the disc begins to cut in deeper than r_{cor}. Its ringshaped domain

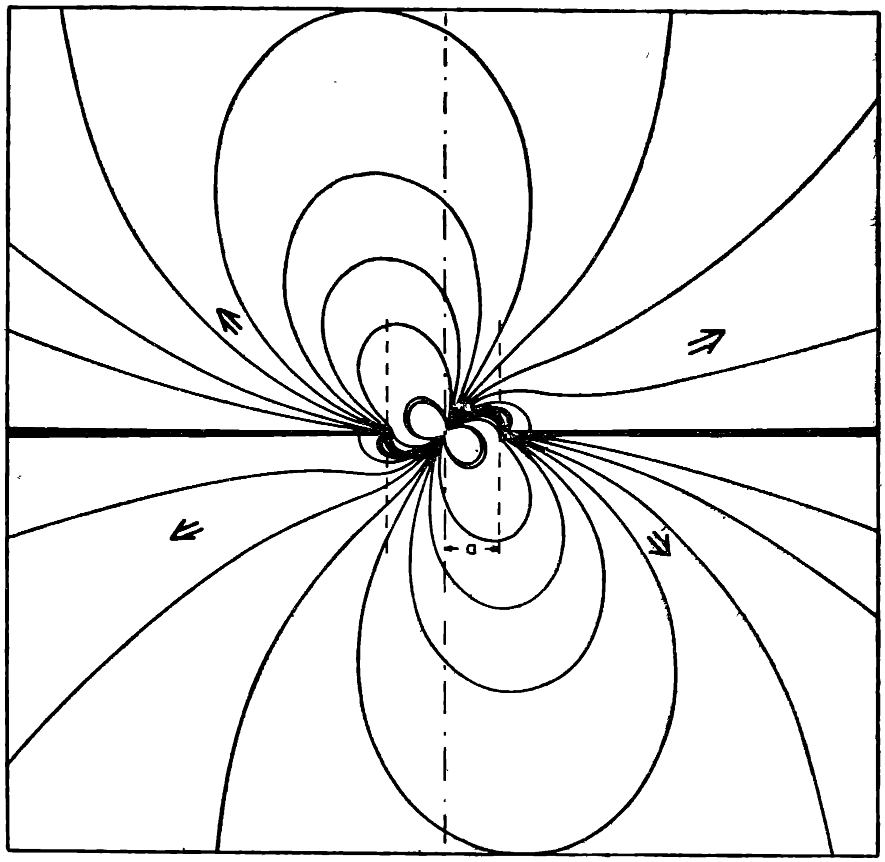

Figure 12. Cut through an inclined (30°) 3-d magnetic dipole confined by a thin, rigid, diamagnetic disc; taken from Kundt & Robnik (1980). In application to the accretion scenario of a spinning, magnetized neutron star, the inner disc radius a is argued to fall just inside the corotation radius r_{cor}. Mass accretion is expected to proceed from very near the inner edge, via field-line-crossing of clumps of matter until (at least) the magnetosphere has an inward slope. If a small fraction of the accreted matter is expelled explosively, it will follow the field lines as indicated by the arrows until magnetic pressures fall below equipartition values; alternatively, a pulsar wind may follow these 'rails'. In reality, the inner disc may be warped and clumpy, and enveloped by a corona.

inside r_{cor} rotates faster than the central star. This means that field lines interacting with the inner edge exert an accelerating torque on the star, so that both accretion and friction at the disc tend to accelerate it. Such a runaway spinup can only be halted by an increasing radiation torque (pulsar torque) of the fast-spinning star. It is possible that the 69 ms x-ray pulsar A 0538–66 realizes this steady state.

In all other cases, the neutron star's spin rate is too low for the radiation torque to be important. Quite likely, accretion near r_{cor} is fast enough to prevent the disc from spreading further in. This interpretation is supported by the fact that we observe erratic fluctuations of the spin period around some equilibrium value rather than a quasi-monotonic decrease. I shall therefore assume that accretion takes place from $a = r_{cor}$ for all rotators except the very fastest.

This assumption differs from the one made by Ghosh & Lamb (1979), *viz.* that field lines can reconnect vertically through the disc (at the local Alfven speed) rather than working their way out radially from the inner edge. Such reconnection would temporarily raise the magnetic tensions, *i.e.* should be energetically forbidden. With this assumption, Ghosh & Lamb estimate field line penetrations into the disc out to more than 30 times the inner radius, and their accretion radius a comes out much smaller than r_{cor}, 0.5 r_{cor} being an upper limit. As stated above, I shall assume $a = r_{cor}$.

The torque T between an accreting neutron star and its feeding disc is the sum of a material and a magnetic torque :

$$T = T_{mat} - T_{mag}. \quad ...(37)$$

T_{mag} can be decomposed into an interaction torque with the disc, T_{disc}, via the field lines intersecting it, and a pulsar-like torque T_{wind} exerted along the 'open' magnetic field lines :

$$T_{mag} = T_{disc} + T_{wind}. \quad ...(38)$$

Here T_{disc} can have either sign, depending on the average axial distance a at which field lines cut through the disc; it is positive (braking) for $a > r_{cor}$. T_{wind} takes care of (i) a possible ejection of matter from the polar caps, (ii) a scratching of matter off the faces of the disc, and (iii) an emission of strong magnetic dipole waves; it is unlikely to be important for spin periods in excess of one second because it is exerted via field lines that bend at much larger radii than that of the inner edge. We have (with equation (36))

$$T_{mat} \approx \dot{M} r_{cor}^2 \Omega = 4 \cdot 10^{34} \text{ erg s}^{-1} \dot{M}_{17} \Omega_0^{-1/3}, \quad ...(39)$$

$$T_{disc} = \frac{1}{4\pi} \iint B_{r\perp} B_{\varphi} r_{\perp}^2 \, d\varphi dz = \pm \xi B^2 r_{cor}^3/4$$

$$= \pm \xi \mu^2 \Omega^2/GM = \pm 5 \cdot 10^{35} \text{ erg } \xi \mu_{31}^2 \Omega_0^2, \quad ...(40)$$

$$T_{wind} \approx 2 \mu_{\perp}^2 \Omega^3/3 c^3 = 2 \cdot 10^{33} \text{ erg } \mu_{\perp,31}^2 \Omega_1^3 \quad ...(41)$$

for an accretion rate $\dot{M}$ in units of 10^{17} g s^{-1} $\approx 10^{-9} M_{\odot}$/yr, a magnetic dipole moment μ in units of 10^{31} G cm^3, and a 'coupling efficiency' $\xi > 0$ which can be both $\ll 1$, for $a \lesssim r_{cor}$, and $\gg 1$ for $a \gtrsim r_{cor}$ when a winding of magnetic field lines by the Keplerian disc generates a strong toroidal magnetic field. T_{disc} is positive (braking) for $a_{eff} > r_{cor}$ and negative for $a_{eff} < r_{cor}$. T_{wind} will subsequently be ignored for $\Omega \lesssim 10$ s^{-1} ($\Omega_1 \lesssim 1$). An accretion rate $\dot{M}$ corresponds to an accretion power

$$L = GM\dot{M}/R = 2 \cdot 10^{37} \text{ erg s}^{-1} \dot{M}_{17} \quad ...(42)$$

for a 1.4 $M_{\odot}$ object.

Now we remember that the observed spin periods are highly variable, implying an (almost) average balance between T_{mat} and T_{disc} (for $\Omega_1 \lesssim 1$). A balance is equally suggested by the long average spinup timescales $P/\langle -\dot{P} \rangle$ compared with the spinup times τ under pure accretion :

$$\tau = I\Omega/T_{mat} = 2 \cdot 10^3 \text{ yr } I_{45} \Omega_0^{4/3}/L_{37}, \quad ...(43)$$

compare figure 13. It implies

$$1 \approx T_{\text{disc}}/T_{\text{mat}} = \xi/Q \qquad \text{...(44)}$$

where from equations (39), (40) and (36)

$$Q := LR(GM)^{2/3}\,\Omega^{-7/3}\,\mu^{-2} = (L\Omega^{-7/3})_{38\cdot 5}\,\mu_{31}^{-2}.$$

Table 2 shows that for $\mu_{31} = 1$, the quantity Q varies from very small values for the fast sources ($0\cdot 8 \cdot 10^{-3}$) to very large values ($3 \cdot 10^{7}$) for the slow sources. This variation, if interpreted as due to a varying magnetic dipole moment μ, would ask for a much wider spread than observed for the radio pulsars, *cf.* figure 3, and in particular for much higher extreme values of μ ($\gtrsim 10^{32\cdot 5}$ G cm^3). I find it more likely that this large range in Q reflects a range in compensation ($\xi \ll 1$) or flux winding ($\xi \gg 1$) respectively, whereby even the highest values of Q correspond to magnetic pressures $B^2/8\pi \lesssim 10^3$ dyn cm^{-2} which are smaller than the expected ram pressures $\rho v^2/2$ near the inner edge of the disc, *i.e.* which can be dynamically supported by the revolving disc material. With this interpretation, the magnetic dipole moments of the x-ray pulsars are comparable to those of the young radio pulsars.

Figure 13 compares the observed average spinup times $P/\langle -\dot{P}\rangle$ with the shortest ones, called τ, permitted by pure accretion (equation (43)) which are drawn in as the lower straight line. It shows that on average, spinup is slower than without disc friction by factors between 3 and $\gtrsim 30$, with no obvious correlation with the scaled (x-ray) luminosity $L\Omega^{-4/3}$. On the other hand, spinup is faster than possible for a degenerate white dwarf with the same luminosity, as is shown by the upper straight line. Figure 13 reassures us that our assumption of $a = r_{\text{cor}}$ is not unreasonable.

Figure 13 does not contain the fastest (69 ms) x-ray pulsar which would be off the frame even if its $\dot{P}$ were negative. For this neutron star, accretion plus disc friction may be balanced by the pulsar-type wind torque. Equations (39) and (41) yield the approximate condition

$$\Omega_{\text{eq}} \approx (3\,c^3\,(GM)^{2/3}\,\dot{M}/2\,\mu_\perp^{2})^{3/10} = 10^{1\cdot 9}\ \text{s}^{-1}\,(\dot{M}_{19}/\mu_{\perp,31}^{2})^{3/10}, \qquad \text{...(45)}$$

corresponding to an equilibrium period $P_{\text{eq}} \approx 2\pi/\Omega_{\text{eq}} = 79$ ms which is close to the one observed. In this model, therefore, the accretion disc of A 0538–66 cuts in deeply beyond r_{cor} such that accretion plus disc friction are balanced by the emission of pulsar waves.

It remains to be discussed why the x-ray pulsars show evolving spinup rates. Our condition (44) implies that the quantity $L/\Omega^{7/3}\xi$ remains constant, or, for an approximately constant ξ, the quantity $\dot{M}/\Omega^{7/3}$. We thus arrive at Savonije's (1978) interpretation that the timescale t for average spinup is related to the evolutionary timescale for binary mass transfer

$$t = \Omega/\langle\dot{\Omega}\rangle = (3/7)\,\dot{M}/\ddot{M}. \qquad \text{...(46)}$$

5.3. *The flickering x-ray sources*

The class of flickering x-ray sources is larger than that of the pulsing sources and may well consist of binary neutron stars with spin rates that have not been detectable. It contains probable neutron-star binaries such as Sco X-1 and 3 U 1700–377, but it also contains bizarre sources such as the low-mass x-ray binary GX 339–4 which shows even quasi-periodic optical intensity fluctuations of order unity, at

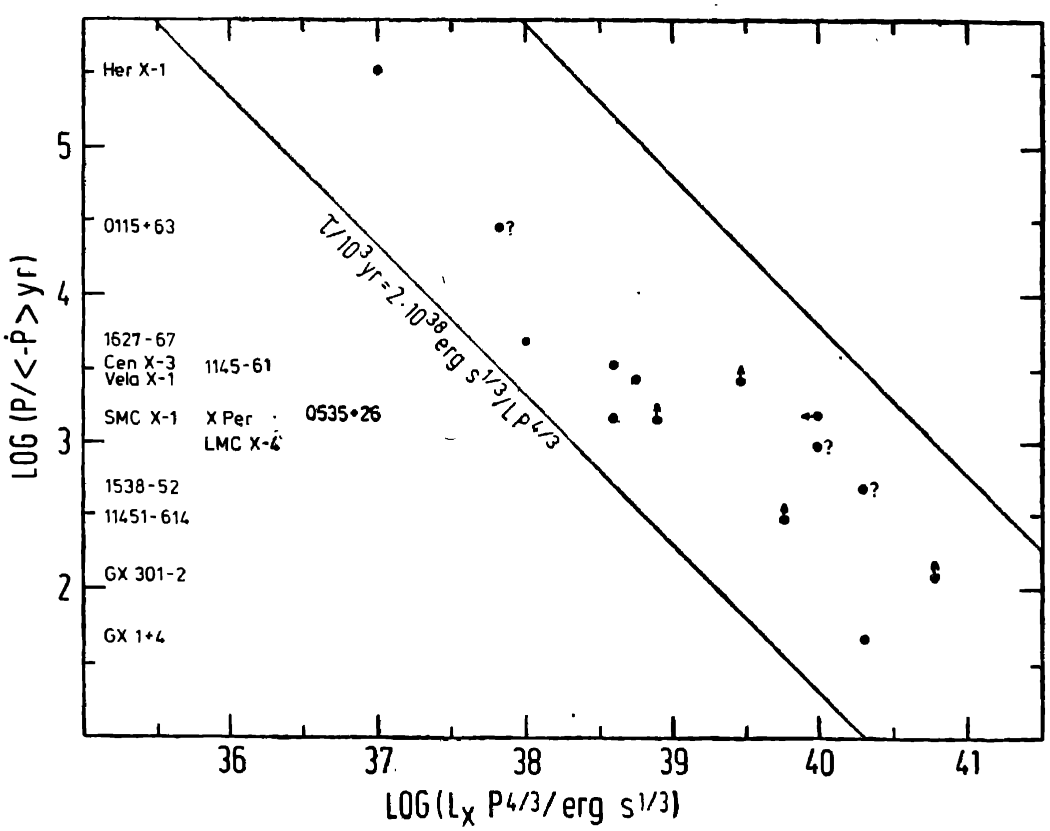

Figure 13. Average spinup timescales $\tau := P/\langle -\dot{P} \rangle$ *versus* $L_X P^{4/3}$ for 14 x-ray pulsars, with L_X = x-ray luminosity; P = pulse period. Plotted are sort of minimal values of τ for several months' averages, as given in table 2. The theoretical minimum of τ (for pure accretion) reads $\tau = 2 \cdot 10^3$ yr $(10^{38}$ erg $s^{1/3}/LP^{4/3})$ and is drawn in, together with the corresponding minimum for a degenerate dwarf; *cf.* equation (43).

repetition times as short as 20 ms (*e.g.* van Paradijs 1983). It even contains the class of blackhole candidates, like Cyg X-1 and LMC X-3.

The x-ray binary Sco X-1 is well-known as a Galactic triple radio source, similar to the extended extragalactic radio sources which are powered by an active galactic nucleus. Its outer radio lobes are probably powered by two relativistic jets (Kundt & Gopal-Krishna 1984), perhaps even by extremely relativistic electron-positron jets (Kundt 1982a). The flickering of Sco X-1 (on timescales from minutes down to less than a second) is likely to be due to magnetobremsstrahlung of the relativistic $e^{\pm}$ which are blown out by the fast spinning neutron star, in the form of a pulsar wind confined by the thermal wind of the companion. This interpretation gives a hint as to how an enormous (optical or x-ray) luminosity can be modulated on timescales which, in the case of GX 339–4, are shorter than the light-travel time needed to illuminate a sufficiently large radiating area with a beam from the centre. The occurrence of beamed extremely relativistic charges in neutron-star binaries is equally indicated by the ultrahard γ-rays ($E \lesssim 2 \cdot 10^{16}$ eV) received during a short ($\approx 1\%$) orbital phase interval from Cyg X-3 (Samorski & Stamm 1983; Dowthwaite *et al.* 1983);

also $3 \cdot 10^{15}$ eV from Vela X-1 (Protheroe *et al.* 1984), $\lesssim 10^{15}$ eV from Her X-1 (*cf.* Dowthwaite *et al.* 1984), and $\approx 10^{16}$ eV rrom LMC X-4 (Protheroe & Clay 1985) An alternative way of UHE γ-ray production is, however, via clumped accretion (Kundt *et al.* 1985).

Intensity flickering is common to both neutron-star sources and blackhole candidates. If such flickering can be produced by extremely relativistic $e^{\pm}$-beams ejected by a fast spinning magnet, we may wonder whether the blackhole candidates do not also contain a fast neutron star. If they did, the large unseen mass in the system ($\gtrsim 5\ M_{\odot}$) would have to be present in a not-too-hot state, perhaps in the form of a massive, self-gravitating accretion disc (Kundt 1979). Such a massive disc would form during an epoch of rapid mass transfer from the companion, instead of the widely preferred common-envelope system (van den Heuvel 1983) which is not easy to form around a neutron star.

5.4. *The bursting sources*

The bursting x-ray sources are bright ($\gtrsim 10^{34}$ erg s^{-1}, mostly $\gtrsim 10^{36}$ erg s^{-1}) x-ray sources which show quasi-regular bursts, at intervals between hours and days (for type I events), of duration less than a minute, with an initial blackbody temperature of $\lesssim 3 \cdot 10^{7}$ K which declines during the burst (Lewin & Joss 1983). The time-averaged burst radiation amouts to about 1% of the steady radiation. The x-ray bursters belong to the class of Galactic bulge x-ray sources which contributes some 80% to the Galactic x-ray power. Their mass-shedding companion tends to be a low-mass star, $M \lesssim 1\ M_{\odot}$, which is small in size so that occultations are rare. Because of the long evolutionary timescale of the companion, these systems are held to be long-lived ($\lesssim 10^{10}$ yr) and are correspondingly rare systems. Bursting and pulsing are (almost) mutually exclusive properties, perhaps controlled by the strength of the magnetic field or by the mode of accretion.

For some time, the absence of occultations in Galactic bulge x-ray sources has been interpreted as due to a thick (in angle) accretion disc which shields the donor star against the x ray point source. This interpretation has recently encountered a number of difficulties (Verbunt *et al.* 1984) and must probably be dismissed: In particular, sharp x-ray eclipses have been discovered in MXB 1659–29, and there are several indications of (pulsed) x-ray reprocessing into optical radiation on the surface of the donor star. Instead, there may be a stalled boundary layer between the windzones of the two stars, of size $\gtrsim l$ sec, and/or a strong wind from the disc which is illuminated during outbursts and gives rise to the (somewhat delayed) optical bursts. An extreme case for this interpretation is SS 433 (Kundt 1981c).

The bursts of type I, whose temperature declines after a sudden initial rise whilst the luminous area remains constant, are thought to be produced by the combustion of matter accreted onto the surface of a neutron star, in the form of a thermonuclear flash. In contrast, the so-called rapid burster shows at the same time bursts of type II, of durations between seconds and hours, with repetition times occasionally as short as several seconds (within the two to six 'active' weeks of a six-months cycle), during which the temperature stays constant ($1.8 \cdot 10^{7}$ K) whilst the intensity (area) decreases. Bursts of type II are thought to be due to discontinuous (spasmodic) accretion onto the whole surface. Important for our understanding of neutron stars is the fact that for both types of bursts, there is a characteristic luminous area πR^2 of radius $R \approx 10$ km. Bursts can therefore measure the radius of a neutron star!

It would be nice to know whether or not the burst sources have lower magnetic fields than young neutron stars. One could argue that the accreting matter would not spread over the whole stellar surface if the magnetic dipole moment was high ($B \gtrsim 10^{10}$ G). Yet accretion could take place in the form of clumps which can penetrate deeply into the magnetosphere and thereby reach large parts of the surface.

An indication that even very old neutron stars can have strong magnetic surface fields comes from the γ-ray burst sources. γ-ray bursts are rare events, with photon energies between 30 keV and many 10^2 keV, total frequency several per month, durations of $\lesssim 0.1$ s, 1 s, or $\gtrsim 10$ s, with several repetitions, and with an isotropic distribution over the sky. Their distribution in number and burst energy is consistent with a nearby ($d \lesssim 10^2$ pc) population of neutron stars radiating near or below their Eddington limit (Mazets *et al.* 1978; Norris *et al.* 1984). The famous γ-ray event of 1979 March 5 has meanwhile recurred at least 16 times and may well be one of the nearest hard x-ray bursters (Golenetskii *et al.* 1984). During the late stage of its first outburst, it showed an 8 s period. In the meantime, three further periodic γ-ray bursts have been recorded, with periods between 4 s and 10 s as expected for dead pulsars (Lipunov *et al.* 1982).

Many spectra of γ-ray bursters show features between 30 and 70 keV which look similar to J. Trümper's cyclotron line in the spectrum of the x-ray pulsar Her X-1, and which have been interpreted as cyclotron emission and/or absorption lines in a magnetic field of several 10^{12} G (Mazets *et al.* 1981). Unless this interpretation is premature, it has provided us with an independent measurement of the (strong) surface magnetic fields of old neutron stars.

Returning once more to the x-ray bursters, we have the yet unsettled question of their origin. If neutron stars derive from massive binaries, how did the present low-mass companion of the neutron star come into existence? X-ray bursters are frequent in globular clusters, several 10^2 times more frequent than in the bulge of the Galaxy, and are therefore often thought to have formed by tidal capture from single or double stars, in the high density core of the globular cluster. There is, however, the problem that one should find correspondingly many cataclysmic variables in globular clusters, more than are observed.

In any case, there remains the problem of how the corresponding Galactic bulge systems have formed, a problem which is aggravated by the fact that the low-mass companion tends to look rather undeveloped (hydrogen-rich). Has a white dwarf been turned into a neutron star by excessive accretion? Arguments against this interpretation have been presented in section 2.1. In spite of strong reservations by friends of mine, I find it plausible that an undeveloped, low-mass companion to a neutron star may form from a massive accretion disc via gravitational instability, when the donor star turns into a supernova. In this scenario, the (rare) low-mass x-ray system would derive from a Cyg X-1-type system after its second supernova explosion.

5.5. SS 433

The Galactic x-ray and radio binary star SS 433 ranges among the most fascinating astrophysical objects of this century, mostly because of its sinusoidally moving spectral lines (of H and He) and relativistic twin jets, and also because of its obvious association with the old supernova remnant W 50. Its kinematic properties are believed to be well understood, such as its distance $d = (5.5 \pm 1)$ kpc, orbital inclination $i = 79°$, beam precession angle $\theta = 20°$, ejection velocity v $= 0.26\,c$, and the geometry of its accretion disc (*e.g.* Margon 1984), but a number of details remain unexplained or even inconsistent (Kundt 1985a). A more satisfactory description of the system can be obtained for $d = (3 \pm 0.5)$ kpc, $i = 67°$, $\theta = 53°$, v $= 0.9999\,c$, and with no signature of an accretion disc in the light curve so far (Kundt 1981c, 1985a). In view of the fact that SS 433 is considered a miniature QSO which should help us understand the more distant extragalactic sources, this ambiguity in the interpretation ought to be soon sorted out.

5.6. The binary pulsars

Table 1 contains the four currently known binary radio pulsars, of which the binary pulsar 1913+16 is the first discovered and most interesting object because it contains essentially two point masses in close, eccentric orbit around each other. One can therefore study a number of general-relativistic effects, among them the emission of gravitational waves and the delay of a signal on passage near a massive object, both of which conform with the predictions of Einstein's theory (see Weisberg & Taylor 1984). Its unseen companionis likely to be another neutron star, and both have masses of $(1.41 \pm 0.04)\,M_\odot$.

How was the binary pulsar formed? From equation (34) with $\epsilon = 0.617$ we may conclude that only 1.6 $M_\odot$ were ejected in the second supernova explosion. At that time, the binary separation was only $(1 - \epsilon)\,a \approx 1$ solar radius; *i.e.* the first-born neutron star was grazing the surface of a rather compact 3 $M_\odot$ star. Tidal interactions should have forced this progenitor star into synchronous rotation and thereby reduced the action of the magnetic dynamo in its core to a very low value (Levy & Rose 1974). We can thus understand the coincidence of a low magnetic dipole moment with the property of still being bound to its elder brother (Kundt 1980).

Straightforward as it sounds, this interpretation is not generally adopted; instead, one attributes the low magnetic field to old age, and postulates that the younger pulsar is invisible because of unfavourable beaming. As discussed in sections 2.3 and 3.1, a small beaming factor causes problems with the already large number of pulsars, and magnetic field decay is open to doubt.

If, on the other hand, closeness of the presupernova companion reduces simultaneously the magnetic dynamo and the binary disruption probability, we can understand the 6 ms pulsar as the product of an even more extreme event (discovery : Boriakoff *et al.* 1983; interpretation : Paczynksi 1983, Joss & Rappaport 1983, Savonije 1983). In this case, the mass of the donor star (with its present $0.3 M_\odot$) has shrunk below the critical value for a supernova explosion. This explanation takes care of two of the three exceptional 'points' in figure 2, below the strip of the other almost 300 pulsars.

There remains, however, the 1.56 ms pulsar with the extremely high spin rate and extremely low $\mu_\perp$ which is unlikely to be very old because it is located very near the Galactic disc (Backer *et al.* 1983; Ashworth *et al.* 1983; Stinebring & Cordes 1983; Becker & Helfand 1983). Several authors have suggested that it received its high spin rate through accretional spinup and that the donor star has subsequently exploded as a supernova or been tidally disrupted (Alpar *et al.* 1982; Fabian *et al.* 1983; Ruderman & Shaham 1983). But spinup via accretion to periods of order ms faces a number of difficulties : To begin with, a massive progenitor system would not live long enough to transfer the needed $\gtrsim 0.1\, M_\odot$ to the neutron star because of the limiting Eddington rate of $1.5 \cdot 10^{-8}\, M_\odot\, \mathrm{yr}^{-1}$ (Arons 1983). A low-mass progenitor system, on the other hand, would be expected to leave a Roche lobe-filling remnant of mass $\approx 10^{-2}\, M_\odot$ which would have been discovered in the timing data. Also, during the long transfer time of $\gtrsim 10^7$ yr, the system would be expected to have moved away from the Galactic disc.

Moreover, a (high) spinrate approaching the ms-range leads to a deformation of the star by centrifugal forces which raises its gravitational binding energy, *i.e.* raises the required energy above that of a rigid rotator (Cowsik *et al.* 1983). At the same time, spinup via accretion gets inefficient when the corotation radius approaches the stellar radius, *cf.* equation (43). Both facts together imply that the transferred mass may have to exceed $0.3\, M_\odot$, which makes the above problems even more severe.

As an alternative possibility, Henrichs & van den Heuvel (1983) propose that the 1.6 ms pulsar has formed through coalescence from a close neutron-star binary in which the large orbital angular momentum is converted into spin angular momentum of the fused object. Perhaps a pulsar as exceptional as this one needs an exceptional explanation of its evolutionary history.

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References

Abbott, D. C., Bieging, J. H. & Churchwell, E. (1984) *Ap. J.* **280**, 671.
Akhiezer, A. I. & Polovin, R. V. (1956) *Sov. Phys. JETP* **3**, 696.
Alpar, M. A., Anderson, P. W., Pines, D. & Shaham, J. (1984a, b) *Ap. J.* **276**, 325; **278**, 791.
Alpar, M. A., Cheng, A.F., Ruderman, M. A. & Shaham, J. (1982) *Nature* **300**, 728.
Alpar, M. A., Langer S. A. & Sauls, J. A. (1984c) *Ap. J.* **282**, 533.
Anderson, B. & Lyne, A. G. (1983) *Nature* **303**, 597.
Arons, J. (1981, 1983) *IAU Symp. No. 95*, p. 69; *Nature* **302**, 301.
Arons, J. & Scharlemann, E. T. (1979) *Ap. J.* **231**, 854.
Ashworth, M., Lyne, A. G. & Smith, F. G. (1983) *Nature* **301**, 313.
Asséo, E., Kennel, C. F. & Pellat, R. (1978) *Astr. Ap.* **65**, 401.
Asséo, E., Llobet, X. & Schmidt, G. (1980) *Phys. Rev.* A **22**, 1293.
Backer, D. C. (1975) *Astr. Ap.* **43**, 395.
Backer, D. C., Kulkarni, S. R. & Taylor, J. H. (1983) *Nature* **301**, 314.
Backus, P. R., Taylor, J. H. & Damashek, M. (1982) *Ap. J. (Lett.)* **255**, L63.
Bartel, N., Sieber, W. & Wolszczan, A. (1980) *Astr. Ap.* **90**, 58.

Baym, G. & Pethick, C. (1979) *A. Rev. Astr. Ap.* **17**, 415.
Becker, R. H. & Helfand, D. C. (1983) *Nature* **302**, 688.
Bignami, G. F., Caraveo, P. A. & Paul, J. A. (1984) *Nature* **310**, 464.
Bisnovatyi-Kogan, G. S., Popov, Yu. P. & Samochin, A. A. (1976) *Ap. Sp. Sci.* **41**, 287.
Blandford, R. D. (1972) *Astr. Ap.* **20**, 135.
Blandford, R. D., Applegate, J. H. & Hernquist, L. (1983) *M. N. R. A. S.* **204**, 1025.
Bodenheimer, P. (1981) *IAU Symp. No. 93*, p. 5.
Boerner, G. (1980) *Phys. Rep.* **60**, 151.
Boriakoff, V., Buccheri, R. & Fauci, F. (1983) *Nature* **304**, 417.
Bradt, H. V. D. & Mc Clintock, E. (1983) *A. Rev. Astr. Ap.* **21**, 13.
Burnard, D. J., Lea, S. M. & Arons, J. (1983) *Ap. J.* **266**, 175.
Che, A. & Kegel, W. H. (1980) *Astr. Ap.* **92**, 204.
Cheng, A. F. & Ruderman, M.A. (1980) *Ap. J.* **235**, 576.
Cordes, J. M., Weisberg, J. M. & Boriakoff, V. (1983) *Ap. J*, **268**, 370.
Cowsik, R., Ghosh, P. & Melvin, M. A. (1983) *Nature* **303**, 308.
Davidson, K. & Fesen, R. A. (1985) *A. Rev. Astr. Ap.*
Demiański, M. & Prószyński, M. (1979) *Nature* **282**, 383.
Deutsch, A. (1955) *Ann. d' Ap.* **18**, 1.
Downs, G. S. (1981) *Ap. J.* **249**, 687.
Dowthwaite, J. C. *et al.* (1983) *Astr. Ap.* **126**, 1.
Dowthwaite, J. C. *et al.* (1984) *Nature* **309**, 691.
Fabian, A. C., Pringle, J. E., Verbunt, F. & Wade, R. A. (1983) *Nature* **301**, 222.
Flowers, E. & Ruderman, M. A. (1977) *Ap. J.* **215**, 302.
Ghosh, P. & Lamb, F. K. (1979a, b) *Ap. J.* **232**, 259; **234**, 296.
Goldreich, P. & Julian, W. H. (1969) *Ap. J.* **157**, 869.
Goldreich, P. & Julian, W. H. (1970) *Ap. J.* **160**, 971.
Golenetskii, S. V., Ilyinskii, V. N. & Mazets, E. P. (1984) *Nature* **307**, 41.
Graser, U. & Schönfelder, V. (1982, 1983) *Ap. J.* **263**, 677; **273**, 681.
Gull, T. R. & Fesen, R. A. (1982) *Ap. J.* (*Lett.*) **260**, L75.
Gullahorn, G. E. & Rankin, J. M. (1982) *Ap. J.* **260**, 520.
Gupta, S. K., Ramana Murthy, P. V., Sreekantan, B. V. & Tonwar, S. C. (1978) *Ap. J.* **221**, 268.
Hankins, T. & Boriakoff, V. (1978) *Nature* **276**, 45.
Hayakawa, S. (1985) *Physics Reports* **121**, 317.
Helfand, D. J. (1983) *IAU Symp. No. 101*, p 471.
Helfand, D. J. & Becker, R. H. (1984) *Nature* **307**, 215.
Henrichs, H. & van den Heuvel, E. P. J. (1983) *Nature* **303**, 213.
Hillebrandt, W. (1984) *Ann. N. Y. Acad. Sci.* **422**, 197.
Holloway, N. J. (1975) *M. N. R. A. S.* **171**, 619.
Holloway, N. J., Kundt, W. & Wang, Y. M. (1978) *Astr. Ap.* **70**, L23.
Howe, S. K. *et al.* (1983) *Ap. J.* **272**, 678.
Isern, J., Labay, J. & Canal, R. (1984) *Nature* **309**, 431.
Izvekova, V. A., Kuzmin, A. D., Malofeev, V. M. & Shitov, Yu. P. (1981) *Ap. Sp. Sci.* **78**, 45.
Jones, P. B. (1980) *Ap. J.* **236**, 661.
Joss, P. C. & Rappaport, S. A., (1983, 1984) *Nature* **304**, 419; *A. Rev. Astr. Ap.* **22**, 537.
Julian, W. H. (1973) *Ap. J.* **183**, 967.
Källman, C. G. (1979) *Fund. Cosmic Phys.* **4**, 167.
Kamper, K. & van den Bergh. S. (1976) *Ap. J.* **32**, 351.
Katz, J. I. (1983) *Astr. Ap.* **128**, L1.
Kennel, C. F. & Coroniti, F. V. (1983) *Ap. J.* **283**, 694.
Krotscheck, E., Kundt, W. & Heintzmann, H. (1975) *Neutron Star Corequakes and Crustquakes II-Strain in the Crust*, received by Ann. of Physics on 9 May.
Kulsrud, R. M., Ostriker, J. P. & Gunn, J. E. (1972) *Phys. Rev. Lett.* **28**, 636.
Kundt, W. (1976, 1977) *Nature* **261**, 673; *Naturwissenschaften* **64**, 493.

Kundt, W. (1979, 1980). *Astr. Ap.* **80**, L7; *Nature* **284**, 246.
Kundt, W. (1982a, b) *Astr. Ap.* **98**, 207; *IAU Symp. No. 95*, p 57.
Kundt, W. (1981c) *Vistas Astr.* **25**, 153.
Kundt, W. (1982a, b) *IAU Symp. No. 97*, p. 265; *Ap. Sp. Sci.* **90**, 59.
Kundt, W., (1983, 1984) *Astr. Ap.* **121**, L15; *Ap. Sp. Sci.* **98**, 275.
Kundt, W. (1985a) SS 433 revisited, *Astr. Ap.*
Kundt, W. (1985b) Supernova Shell Structure, in : *Fairfax meeting on Crab Nebula and Related SNRs* (Oct. 84), (ed : M. Kafatos) to appear in Cambridge Univ. Press.
Kundt, W. (1985c) *Particle acceleration by pulsars*, in "*Erice Lecture*" (Nov. 1984) to appear in Reidel (ed.: M. M. Shapiro).
Kundt, W. & Gopal-Krishna (1984) *Astr. Ap.* **136**, 167.
Kundt, W. & Krotscheck, E. (1980) *Astr. Ap.* **83**, 1.
Kundt, W. Özel, M. E. & Ercan, E. N. (1985) Are the galactic-bulge x-ray sources magnetized? submitted to *Astr. Ap.*
Kundt, W. & Robnik, M. (1980) *Astr. Ap.* **91**, 305.
Levy, E. H. & Rose, W. K. (1974) *Nature* **250**, 40.
Lewin, W. H. G. & Joss, P. C. (1983) in *Accretion-Driven Stellar X-ray Sources* eds : W. H. G. Lewin & E. P. J. vanden Heuvel) Cambridge Univ. Press, p. 41.
Lipunov, V. M. (1982) *Ap. Sp. Sci.* **82**, 343.
Lipunov, V. M., Moskalenko, E. I. & Shakura, N. I. (1982) *Ap. Sp. Sci.* **85**, 459.
Lohsen, E. H. G. (1981) *Ap. Ap. Suppl.* **44**, 1.
Lyne, A. G. (1982) in : *Supernovae, A Survey of Current Research* (eds : M. J. Rees & R. J. Stoneham) Reidel, p. 405.
Lyne, A. G., Anderson, B. & Salter, M. J. (1982) *M. N. R. A. S.* **201**, 503.
Manchester, R. N., Durdin, J. M. & Newton, L. M. (1985) *Nature* **313**, 374.
Manchester, R. N. & Taylor, J. H. (1977) *Pulsars*, Freeman.
Manchester, R. N. & Taylor, J. H. (1981) *Astr. J.* **86**, 1953.
Margon, B. (1984) *A. Rev. Astr. Ap.* **22**, 507.
Marsden, P. L. *et al.* (1984) *Ap. J. (Lett.)*, **278**, L29.
Mazets, E. P., Golenetskii, S. V., Aptekar', R. L., Gur'yan, Yu. A. & Il'inskii, V. N. (1981) *Nature* **290**, 378.
McCulloch, P. M., Hamilton, P. A., Royle, G. W. R. & Manchester, R. N. (1983) *Nature* **302**, 319.
Melrose, D. B. (1980) *Plasma Astrophysics*, Gordon & Breach.
Mestel, L. (1971) *Nature Phys. Sci.* **233**, 149.
Michel, F. C. (1982) *Rev. Mod. Phys.* **54**, 1.
Michel, F. C. & Dessler, A. J. (1983) *Nature* **303**, 48.
Middleditch, J. & Pennypacker, C., (1995) *Nature* **313**, 659.
Miller, E. W. (1973) *Pub. Astr. Soc. Pacific* **85**, 764.
Moffatt, H. K. (1978) *Magnetic Field Generation in Electrically Conducting Fluids*, Cambridge Univ. Press.
Narayan, R. & Vivekanand, M. (1982, 1983) *Astr. Ap.* **113**, L3; **122**, 45,
Norris, J. P., Cline, T. L., Desai, U. D. & Teegarden, B. J. (1984) *Nature* **308**, 434.
Ochelkov, Yu. P. & Usov, V. V. (1980) *Ap. Sp. Sci.* **69**, 439.
Okamoto, I. (1975) *M. N. R. A. S.* **170**, 81.
Pacini, F. & Salvati, M. (1983) *Ap. J.* **274**, 369.
Paczynski, B. (1983) *Nature* **304**, 421.
Protheroe, R. J. & Clay, R. W. (1985) *Nature* **315**, 205.
Protheroe, R. J., Clay, R. W. & Gerhardy, P. R. (1984) *Ap. J.* **280**, L47.
Radhakrishnan, V. (1982) *Contemp. Phys.* **23**, 207.
Rankin, J. M. (1983) *Ap. J.* **274**, 333; 359.
Rappaport, S. A. & Joss, P. C. (1983) *Accretion-Driven Stellar X-ray Sources* (eds : W. H. G. Lewin & E. P. J. van den Heuvel) Cambridge Univ. Press, p. 1.
Rees, M. J. (1981) *IAU Symp. No. 97*, p. 211.

Ritchings, R. T. (1976) *M.N.R.A.S.* **176**, 249.
Ruderman, M. A. & Shaham, J. (1983) *Nature* **304**, 425.
Ruderman, M. A. & Sutherland, P. G. (1973) *Nature Phys. Sci.* **246**, 93.
Samorski, M. & Stamm, W. (1983) *Ap. J.* 268, L17.
Sanford, P. W., Laskarides, P. & Salton, J. (1982) *Galactic X-ray Sources*, Wiley.
Savonije, G. J. (1978, 1983) *Astr. Ap.* **62**, 317; *Nature* **304**, 422.
Scharlemann, E. T. (1974) *Ap. J.* **193**, 217.
Shapiro, S. L. & Teukolsky, S. A. (1983) *Blackholes, White Dwarfs, and Neutron Stars*, Wiley.
Sieber, W. & Wielebinski, R. (eds.) (1981) *IAU Symp. No. 95.*
Skinner, G. K. *et al.* (1982) *Nature* **297**, 568.
Srinivasan, G. & Dwarakanath, K. S. (1982) *J. Ap. Astr.* **3**, 351.
Stella, L. *et al.* (1985) *Ap. J.* to appear.
Stinebring, D. R. & Cordes, J. M. (1983) *Nature* **306**, 349.
Stone, R· (1979) *Ap. J.* **232**, 520.
Sutherland, P. G. (1979) *Fund. Cosmic Phys.* **4**, 95.
Tammann, G. A. (1977) *Ann. N. Y. Acad. Sci.* **302**, 61.
Tuffs, R. J. (1983) *IAU Symp. No. 101*, p. 49.
van Paradijs, J. (1983) in *Accretion-Driven Stellar x-ray Sources* (eds : W. H. G. Lewin & E. P. J. van den Heuvel), Cambridge Univ. Press, p. 189.
van den Heuvel, E. P. J. (1983) in *Accretion-Driven Stellar X-ray Sources* (eds · W. H. G. Lewin & E. P J. van den Heuvel) Cambridge Univ. Press, p. 303.
van den Heuvel, E. P. J. & Taam, R. E. (1984) *Nature* **309**, 235.
Verbunt, F., Fabian, A. C. & Rees, M. J. (1984) *Nature* **309**, 331.
Vivekanand, M. & Narayan, R. (1981) *J. Ap. Astr.* **2**, 315.
Wassermann, I. & Shapiro, S. L. (1983) *Ap. J.* **265**, 1036.
Wegener, U. (1977) *Diplomarbeit*, Hamburg.
Weiler, K. W., Sramek, R. A., van der Hulst, J. M. & Panagia, N. (1983) *IAU Symp. No. 101*, p. 171.
Weisberg, J. M. & Taylor, J. H. (1984) *Phys. Rev. Lett.* **52**, 1348.
Wilson, D. & Rees, M. (1978) *M. N. R. A. S.* **185**, 297.
Zeldovich, Ya. B., Ruzmaikin, A. A. & Sokoloff, D. D. (1983) *Magnetic Fields in Astrophysics*, Gordon & Breach.