observed in SGCs is typical of a propagating star formation process. Various pieces of evidence in favor of such a process have been noted more than once before (e.g., Ref. 17). The simplest explanation for the existence of an age gradient of star groups in SGCs is a shock wave which propagates through a giant H I cloud and initiates star formation in it.

The age gradient discovered across the Sagittarius—Carina arm indicates that this arm was formed by a spiral density wave. The velocity of propagation of star formation in SGC2 turned out to be 45-100 km/sec. Rough estimates of the rotation rate of the spiral pattern yield $\Omega_p \approx 17$-25 km/sec·kpc, which agrees with the data of Ref. 18 obtained for part of the Sagittarius—Carina arm in quadrant IV. The corresponding corotation radius is ~8.8-13.0 kpc. Then in the investigated parts of the Cygnus or Perseus—Cassiopeia arms, as well as in the S6 arm lying near the corotation radius in the galaxy M31 (Ref. 5 and the references therein), there should be no age gradient across the arm. In fact, in the Cygnus arm (SGC3) the ages of star groups increase along the arm. The investigated part of the Perseus—Cassiopeia arm is probably located outside the corotation radius.

The estimates of the propagation velocity of star formation obtained above were determined using age gradients observed in the associations in SGC2. The same velocities found from clusters inside and outside SGC2 ($\Delta t \approx 2\cdot 10^7$ yr) turned out to be considerably lower, ~10-16 km/sec.

In each complex, star formation propagates in a certain direction (from one edge of the H I supercloud toward the other), independently of neighboring complexes. This supports Efremov's suggestion [2,5] that giant clouds of neutral hydrogen are the starting units for star formation. The situation described above indicates a physical relation between different components of complexes, which were identified earlier only on the basis of close promimity in space.

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Cyclotron instability and the generation of radio emission in pulsar magnetospheres

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It is shown that cyclotron instability can develop in the magnetosphere of a typical pulsar if the magnetic field near its surface differs greatly from a dipole by having a radius of curvature of the field lines of the order of the radius of the neutron star. This instability develops near the light cylinder of the pulsar and leads to the generation of radio emission. The properties of this radio emission are examined.

1. The most advanced models of pulsars [16,23,1,2] are based on the Sturrock hypothesis [18] that the plasma of pulsar magnetospheres consists primarily of electron—positron pairs created by gamma rays in a strong magnetic field. These pairs very quickly lose the momentum component perpendicular to the magnetic field because of synchrontron losses. As a result, the electron—positron plasma, in escaping from the vicinity of the pulsar, is one-dimensional, i.e., the electrons and positrons move essentially along the magnetic field lines. At some distance from the pulsar, the one-dimensional relativistic electron—positron plasma can become unstable to the generation of transverse electromagnetic waves in cyclotron resonance with the high-energy particles.[12,11] The development of this instability,

called the cylotron instability, is due to the one-dimensionality of the plasma and the high asymmetry of the energy distribution function of the electrons and positrons in the pulsar plasma, that is, the presence of an extended high-energy "tail." It has been shown [12,11] that when the external magnetic field of a neutron star has the dipole configuration, the cyclotron instability can develop in the magnetospheres of only a few sufficiently young pulsars (such as PSR 0531 + 21 and PRS 0833 − 45), and then only under the additional condition that the plasma density in the magnetospheres of these pulsars is significantly greater than the value expected within the scope of the standard pulsar models developed in Refs. 16, 23, 1, 2. We note that the assumption of a strict dipole configuration of the external magnetic field of neutron stars, although very frequently used in work on the theory of pulsars, is unlikely for real pulsars. It is most likely that the magnetic field near the surface of a neutron star differs greatly from a dipole and has a radius of curvature R_c of the field lines of the order of the radius of the neutron star, $R \approx 10^6$ cm. [16] This is indicated by observational data on pulsars[7] and accreting neutron stars in binary systems such as Her X-1. [13] It is shown in the present paper that cyclotron instability can develop in the magnetosphere of a typical pulsar with standard parameters of the pulsar plasma if the magnetic field near the pulsar differs greatly from a dipole in having the value of R_c of the order of the radius of the neutron star. Cyclotron instability develops near the light cylinder of pulsars and leads to radio emission with properties close to those observed.

2. According to Refs. 16, 23, 1, and 2, near the magnetic poles of pulsars there are regions with a strong electric field $\mathbf{E}_\parallel = (\mathbf{E}, \mathbf{B})\mathbf{B}/|\mathbf{B}|^2$ along the magnetic field $\mathbf{B}$. These regions are called the polar gaps. Under the action of the electric field $E_\parallel$ in the polar gap, the primary particles (electrons or positrons, depending on the direction of the field $\mathbf{E}_\parallel$) are accelerated to energies corresponding to the Lorentz factor [16]

$$\gamma_b \approx 3\cdot 10^6 \left(\frac{B_0}{10^{12}\ \mathrm{G}}\right)^{-1/7} \left(\frac{P}{1\mathrm{sec}}\right)^{-1/7} \left(\frac{R_c}{R}\right)^{4/7}, \tag{1}$$

where B_0 is the magnetic field intensity at the surface of the pulsar, and P is its rotation period.

The ultrarelativistic ($\gamma_b \gg 1$) primary particles, in moving essentially along the curved magnetic field lines of the pulsar, generate curvature radiation, whose maximum occurs at photons with energy

$$\mathscr{E}_c = \frac{3}{2}\frac{\hbar c}{R}\gamma_b^3 \approx 10^3 \left(\frac{B_0}{10^{12}\ \mathrm{G}}\right)^{-3/7} \left(\frac{P}{1\mathrm{sec}}\right)^{-3/7} \left(\frac{R_c}{R}\right)^{5/7}\ \mathrm{MeV}, \tag{2}$$

i.e., gamma rays.

In the propagation process, a large part of this γ-radiation is dissipated in the magnetic field of the pulsar with the creation of electron−positron pairs ($\gamma + B \to c^+ + c^- + B$). The differential spectrum of the pairs created has the form $dn/d\gamma \propto \gamma^{-2/3}$ for

$$\gamma \lesssim \gamma_t = \frac{\mathscr{E}_c}{2mc^2} \approx 10^3 \left(\frac{B_0}{10^{12}\ \mathrm{G}}\right)^{-3/7} \left(\frac{P}{1\ \mathrm{sec}}\right)^{-3/7} \left(\frac{R_c}{R}\right)^{5/7} \tag{3}$$

and an exponential cut off for $\gamma > \gamma_t$. [19,16]

The created pairs have nonzero values of the pitch angles in the magnetic field $\mathbf{B}$ and generate synchronous γ-radiation which is also dissipated in the magnetic field with the creation of second-generation electron−positron pairs. The differential spectrum of the second-generation electron−positron pairs is a power law ($dn/d\gamma \propto \gamma^{-3/2}$) for $\gamma \lesssim 0.03\gamma_t$ and an exponential law ($dn/d\gamma \propto \exp\cdot\{\gamma/0.03\gamma_t\}$) for $\gamma > 0.03\gamma_t$. [19]

At low energies, the energy spectrum of the escaping particles cuts off sharply for

$$\gamma \lesssim \gamma_p \approx (n-1)\frac{R_c}{R}\left[1+\left(\frac{B_0}{2\cdot 10^{12}\ \mathrm{G}}\right)^2\right]^{-1/2}, \tag{4}$$

due to an exponential increase in the mean free path of low-energy gamma rays in the magnetic field (Eq. (4) is applicable for $\gamma_p > 1$).

It is assumed here and below that the magnetic field in the magnetosphere of the pulsar varies with distance r from the center according to the following law:

$$B(r) \approx \begin{cases} B_0 (R/r)^n & \text{for } R \leqslant r \lesssim \text{several } R, \\ B_0^d (R/r)^3 & \text{for } R \ll r < c/\Omega, \end{cases} \tag{5}$$

where n is a constant ($n > 3$), B_0^d is the dipole component of the magnetic field at the surface of the neutron star, ($B_0^d < B_0$), $\Omega = 2\pi/P$ is the rotational angular velocity of the pulsar, and c/Ω is the radius of the light cylinder of the pulsar.

For the parameters of the magnetic field near the surface of the neutron star which seem to be most likely for most pulsars ($n \approx 4$-5, $R_c \approx R$, $B_0 \approx (1\text{-}3)\cdot 10^{12}$ G), we have from Eq. (4) $\gamma_p \approx 2$-4.

The electrons and positrons escaping from the vicinity of the pulsar along its magnetic field lines can be divided into three components: 1) electron−positron plasma with $\gamma \sim \gamma_p$, 2) a high-energy plasma "tail" with $\gamma \sim \gamma_t$, and 3) a primary beam with $\gamma \sim \gamma_b$.

The number density of primary-beam particles is [6,16,1]

$$n_b \approx \frac{\Omega B}{2\pi c e} \approx 7\cdot 10^{10}\left(\frac{B}{10^{12}\mathrm{G}}\right)\left(\frac{P}{1\mathrm{sec}}\right)^{-1}\ \mathrm{cm}^{-3}. \tag{6}$$

For typical pulsars ($B_0 \sim 10^{12}$ G, $P \sim 1$ sec), the energy is distributed approximately equally among all three components of electrons and positrons. Hence it follows that

$$n_p \approx \frac{\gamma_b}{2\gamma_p} n_b, \quad \frac{n_t}{n_p} \approx \frac{\gamma_p}{\gamma_t}. \tag{7}$$

The total flux of energy carried away from the vicinity of the neutron star by the escaping particles can reach [16]

$$L_\pm \approx 10^{30}\left(\frac{B_0}{10^{12}\ \mathrm{G}}\right)^{6/7}\left(\frac{P}{1\mathrm{sec}}\right)^{-15/7}\ \mathrm{erg/sec}. \tag{8}$$

3. Cyclotron instability of a relativistic one-dimensional electron−positron plasma with an asymmetric distribution function (the presence of a greatly extended high-energy "tail") has been investigated in detail. [10,12,11] It has been shown that the distance from the center of the pulsar to the region where cyclotron instability develops is

$$r_c \approx R\left(\frac{\gamma_p^2}{\gamma_t}\right)^{1/3}\left(\frac{\omega_{B_0}}{\omega_{p_0}}\right)^{2/3}, \tag{9}$$

where

$$\omega_{B_0}=\frac{eB_0^d}{mc}\approx 1.76\cdot 10^{19}\left(\frac{B_0^d}{10^{12}\,\Gamma\,\mathrm{G}}\right)\mathrm{sec}^{-1},$$

$$\omega_{p_0}=\sqrt{\frac{4\pi e^2 n_{p_0}}{m}}\approx 10^{10}\left(\frac{B_0^d}{10^{12}\mathrm{G}}\right)^{1/2}\left(\frac{P}{1\mathrm{sec}}\right)^{-1/2}\left(\frac{\gamma_b}{\gamma_p}\right)^{1/2}\mathrm{sec}^{-1} \tag{10}$$

are the cyclotron and plasma frequencies at the surface of the pulsar.

Expression (9) for r_c is valid only for $r_c \lesssim c/\Omega$. If the value of r_c calculated from Eq. (9) is greater than c/Ω, then cyclotron instability cannot develop in the magnetosphere of the pulsar.[11] In other words, if cyclotron instability is absent from the near zone of the pulsar, then it will also be absent from its wave zone.

The logarithmic increment for cyclotron instability in the pulsar system is

$$\Gamma_c \approx \frac{\pi}{2}\frac{n_t}{n_p}\left(\frac{\omega_p}{\omega_B}\right)^4\frac{\omega_B}{\gamma_p^3}. \tag{11}$$

Cyclotron instability develops in the magnetosphere of a pulsar when the following conditions are satisfied: a) the distance r_c from the center of the neutron star to the region of instability is no greater than the radius of the light cylinder of the pulsar, i.e., $r_c \lesssim c/\Omega$; b) the characteristic time for instability development $\tau_c = (\Gamma_c)^{-1}$ is less than the time for escape of the relativistic plasma beyond the light cylinder of the pulsar $\tau_0 = (c/\Omega)/c = \Omega^{-1}$.

The conditions for cyclotron instability can be rewritten with Eqs. (1), (3), (7), (9)-(11) as

$$\gamma_p \lesssim 7\eta^{4/27}\left(\frac{B_0^d}{10^{12}\,\mathrm{G}}\right)^{-11/21}\left(\frac{P}{1\,\mathrm{sec}}\right)^{10/21}\left(\frac{R_c}{R}\right)^{3/7}, \tag{12}$$

$$\gamma_p \lesssim 9\eta^{3/28}\left(\frac{B_0^d}{10^{12}\,\mathrm{G}}\right)^{3/7}\left(\frac{P}{1\,\mathrm{sec}}\right)^{17/21}\left(\frac{R_c}{R}\right)^{5/28}, \tag{13}$$

where $\eta = B_0^d/B_0 \lesssim 1$.

Conditions (12) and (13) are satisfied for typical pulsars ($P \approx 1$ sec, $B_0^d \approx 10^{12}$ G) with a magnetic field differing significantly from a dipole ($R_c \approx R$, $\gamma_p \approx 2$-4, and $\eta \approx 0.3$), and cyclotron instability can develop in their magnetospheres.

For a dipole configuration of the external magnetic field of a pulsar, for which $R_c \lesssim (4/3)\cdot(c/\Omega R)^{1/2} R \approx 10^2 (P/1\ \mathrm{sec})^{1/2} R$, it follows from (4), (12), and (13) that cyclotron instability does not develop because of the large value of γ_p.

4. The development of cyclotron instability in the magnetosphere of a pulsar leads to the generation of electromagnetic oscillations with a near-vacuum ($\omega \approx kc$).

The wave number k_{res} of the electromagnetic oscillations in cyclotron resonance with the high-energy particles moving with the Lorentz factor γ_{res}($\gamma_p \ll \gamma_{res} \lesssim \gamma_t$) is[11]

$$k_{res} \approx \frac{\omega_B^3\gamma_p^3}{c\gamma_{res}\omega_p^2}. \tag{14}$$

At the distance from the pulsar $\tilde{r}_c$ ($r_c \lesssim \tilde{r}_c \lesssim c/\Omega$) at which the condition $(\omega_p)^2 \approx \gamma_p^2/\gamma_{res}$ is satisfied,[12] electromagnetic oscillations are generated at a frequency

$$\omega \approx k_{res}c \approx \omega_B(r=\tilde{r}_c)\gamma_p. \tag{15}$$

Substituting Eqs. (9) and (10) into (15), we obtain the frequency of the oscillations generated by particles with $\gamma \approx \gamma_{res}$:

$$\nu = \frac{\omega}{2\pi} \approx \frac{\gamma_{res}\gamma_b}{\gamma_p^2}P^{-1}. \tag{16}$$

The radiation generated as cyclotron instability develops in the magnetosphere of a pulsar is concentrated in the frequency interval from ν_{min} to ν_{max}, where

$$\nu_{min} \approx \frac{\omega_B(r=c/\Omega)}{2\pi}\gamma_p \approx 20\gamma_p\left(\frac{B_0^d}{10^{12}\,\mathrm{G}}\right)\left(\frac{P}{1\,\mathrm{sec}}\right)^{-3}\mathrm{MHz}, \tag{17}$$

$$\nu_{max} \approx \frac{\omega_B(r=r_c)}{2\pi}\gamma_p$$
$$\approx \frac{\gamma_t\gamma_b}{\gamma_p^2}P^{-1} \approx 3\cdot 10^3\gamma_p^{-2}\left(\frac{B_0}{10^{12}\,\mathrm{G}}\right)^{-4/7}\left(\frac{P}{1\mathrm{sec}}\right)^{-11/7}\mathrm{MHz}. \tag{18}$$

For typical pulsars ($B_0 \approx 10^{12}$ G, $P \sim 0.5$-1 sec, and $\gamma_p \approx 2$), we have $\nu_{min} \approx 10$-10^2 MHz and $\nu_{max} \sim$ several GHz, i.e., the radiation generated is in the radio range.

The radio spectrum cuts at low frequencies ($\nu \lesssim \nu_{min}$) and turns over at high frequencies ($\nu \gtrsim \nu_{max}$). These properties of pulsar radio emission to be expected in the cyclotron model agree well with observational data (see Ref. 8 and the references therein).

A decrease in the frequency of the radio emission during its escape from the pulsar magnetosphere is possible due to nonlinear processes.[11,5] These nonlinear processes include, for example, decay of a transverse electromagnetic wave (t-wave) into a longitudinal wave (ℓ-wave) and a t-wave with a frequency lower than the original one ($t \to \ell + t'$). As a result of this decay, the frequency of the low-frequency cutoff in the pulsar radio-emission spectrum can decrease to $\sim\nu_p\sqrt{\gamma_p}$ taken at $r \sim c/\Omega$, which is 1-2 orders of magnitude lower than the frequency ν_{min} given by Eq. (17).

According to this model, the high-frequency cutoff must be near ν_{max}, since there are no sufficiently effective mechanisms to pump the photons to higher frequencies.

As a result of quasilinear relaxation, about half the energy in the high-energy "tail" of the electron–positron plasma is transformed into oscillation energy.[10,11] Consequently, the radio luminosity of the pulsar can reach $\sim 0.1L_+$ (see Eqs. (7) and (8)). This value is quite sufficient to explain the observed radio luminosity of pulsars.

The question about the nature of the directionality of the radio emission of pulsars is trivially resolved, if this radiation is generated deep within the light cylinder of the pulsar, i.e., at a distance from the center of the neutron star $r \ll c/\Omega$. It is then usually assumed that the beam pattern of the radio emission is close to the narrow cone formed by the tangents to the open magnetic field lines escaping beyond the light cylinder of the pulsar.

In our model, when the radio emission of a pulsar is generated near its light cylinder and emerges along the tangents to the magnetic field lines, one must assume, to explain the strong directionality of the radio emission, that the radio-emission region has a small dimension $\Delta\ell$ across the magnetic field lines compared to the radius of the light cylinder ($\Delta\ell \lesssim 0.1c/\Omega$).

Some region near the light cylinder may be distinguished either by a significant difference of the parameters of the escaping plasma from the average values at the same distance r, or by a difference of the magnetic field parameters. Let us assume, for example, that the value of γ_p of the escaping plasma varies by just a small factor across the magnetic field lines, that is, let $\gamma_p \approx \gamma_p^{min} = 2$ along some sufficiently narrow magnetic flux tube with transverse dimension $\Delta\ell \lesssim 0.1\ c/\Omega$ at $r \approx c/\Omega$ and $\gamma_p \approx \gamma_p^{max} = 5\text{-}10$ along most of the open field lines. Then the conditions (12) and (13) are satisfied by a wide margin in the region penetrated by the magnetic field lines along which $\gamma_p \approx \gamma_p^{min}$, and cyclotron instability develops readily at $r \sim c/\Omega$. Over most of the light cylinder region, where $\gamma_p \approx \gamma_p^{max}$, the cyclotron instabilty may not develop, but if it does, the density of the signals generated must be significantly lower because of the small number of particles of the plasma "tail" in resonance with them. We note that the maximum frequency of the oscillations ν_{max} generated via cyclotron instability in the vicinity of field lines with $\gamma_p \approx \gamma_p^{max}$ is shifted to lower frequencies by a factor $\sim(\gamma_p^{max}/\gamma_p^{min})^2 \sim 6\text{-}25$ (see Eq. (18) as compared to the frequency of the oscillations generated in the regions with $\gamma_p \approx \gamma_p^{min}$. Consequently, if cyclotron instability develops in regions with $\gamma_p \approx \gamma_p^{max}$ at $r \sim c/\Omega$, weaker quasi-isotropic low-frequency radio emission ($\nu \lesssim 100\text{-}400$ MHz) will then be generated. Since the slope of pulsar radio spectra in the main pulse increases with increasing frequency, the quasi-isotropic component, whose spectrum is similar to the spectrum of the pulsed emission but shifted to lower frequencies, must then have a larger value of the spectral index $\alpha(I_\alpha \propto \nu^{-\alpha})$ than the pulsed emission at the same frequency. The existence of such quasi-isotropic low-frequency pulsar radio emission with a significantly steeper spectrum than in the radio pulses has been reported repeatedly.[3,4,14,17]

Let us now examine the expected frequency dependence of the radio-emission pulse width. As plasma escapes, cyclotron instability begins at first for the most energetic particles of the plasma "tail" ($\gamma_{res} \simeq \gamma_t$). As follows from (16), the highest-frequency radiation is then generated. With recession from the pulsar, particles with lower and lower energies begin to participate in the development of the instability. It is seen from (9) that the distance $\tilde{r}_c$ at which radiation is generated by particles with $\gamma = \gamma_{res} > \gamma_t$ depends on γ_{res} as $\tilde{r}_c \propto \gamma_{res}^{-1/3}$. From this and Eq. (16), we obtain the frequency dependence of the distance $\tilde{r}_c$ at which radiation at frequency ν is generated: $\tilde{r}_c \propto \nu^{-1/3}$. As long as the distortion of the dipole field due to the closeness of the light cylinder is small, we must have the following frequency dependence of the radio pulse width $\Delta\phi$:

$$\Delta\varphi \propto (\Omega\tilde{r}_c/c)^{1/2} \propto \nu^{-1/6} \approx \nu^{-0.17}. \qquad (19)$$

Such a frequency dependence of the average width of the radio pulses expected theoretically is close to that observed from pulsars.[8]

5. The upper boundary of the polar gap in the various models [16,23,1,2] is determined by the same condition: the primary particles, in crossing the gap, are accelerated to energies such that copious electron—positron pair production begins with the absorption of the curvature gamma rays generated by the primary particles. Thus, the Lorentz factor of the primary particles emerging from the gap is almost the same in all the polar-gap models. The energy spectra of the particles of the electron—positron plasma produced near the pulsar with the absorption of γ-radiation must also be nearly the same (for the same geometry of the magnetic field of the neutron star). Consequently, the results obtained above for cyclotron instability and radio emission in the magnetospheres of pulsars is applicable to all pulsar models [16,23,1,2] based on the assumption of the exisdtence of a polar gap.

We note in conclusion that the generation of radio emission in the magnetospheres of pulsars can result not only from the cyclotron instability examined here, but also from other plasma processes, for example, two-stream instability that develops in the presence of strongly time-dependent flow of electron—positron plasma from the vicinity of a neutron star.[21,22,20] The existence of several mechanisms for the generation of pulsar radio emission is suggested by the great differences in the observed properties of different components of radio pulses.[15,9]

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A mechanism for generating energetic particles and producing coronae in x-ray binaries

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The electric fields generated by a differentially rotating accretion disk can accelerate protons to high energies by an electrostatic mechanism. In an x-ray binary system, protons as energetic as 3 TeV could be expelled perpendicular to the symmetry plane of the magnetized disk at an intensity as high as 10^{36} erg/sec. This mechanism might play an important role in forming an x-ray corona around the system.

Now that ultrahigh-energy γ rays ($E_\gamma \sim 10^{12}$-10^{17} eV) have been detected from several x-ray binary sources, attempts are being made [1-3] to review and to update the classical ideas as to the emission mechanisms and nature of the components of such binary systems as Cygnus X-3, Hercules X-1, and Vela X-1. Apart from the optical range, throughout the spectrum the Cyg X-3 source has the most comprehensive experimental information currently available. [4-6]

When the spectral properties of this source are examined, certain empirical results stand out that do not fit into the classical α-model of disk-type accretion. Among these are the radio flares[7] of 1972 and 1988; the nonthermal part of the spectrum, which conforms to an $E^{-\alpha}$ power law with a spectral index $\alpha = 2.5$ in the hard x-ray and soft γ-ray regions; and finally, the high-energy ($E_\gamma \gtrsim 10^{11}$ eV) γ-ray emission.

Another salient peculiarity is the extensive x-ray corona that envelops the Cyg X-3 system. Similar halos have been encountered [8,9] around other x-ray binaries: the sources 4U 2129 + 49 and 4U 1822−37. This phenomenon too may be regarded as a distinctive one, inasmuch as the luminosity of the sources observed falls below the Eddington limit, so that the mechanism of halo formation through supercritical accretion is not applicable.

In this letter we shall retain the classical premises regarding the objects' nature; that is, we shall assume that systems of this kind have a neutron star as one component, while the companion star is overflowing its Roche lobe. The x-ray emission mechanism will be disk accretion onto the neutron star, whose spin period should in this event not be too short.

Our theoretical model for the emission proces will entail magnetized disk-type accretion in which vital roles are played not only by the viscosity of the disk plasma but also by the magnetic field generated through differential rotation of the disk material. The model will be treated in the MHD approximation.

1. We begin by describing the magnetized-disk model. The matter accreted by the neutron star will posses the companion star's poloidal magnetic field, expressed in a cylindrical coordinate system as $(B_r, 0, B_z)$. As the material enters the differentially rotating disk region, the randomly oriented poloidal field $\mathbf{B}_p$ of the plasma will become redistributed into a disk toroidal field $\mathbf{B}_t$ and will be pumped by the dynamo mechanism, producing a mixed field $\mathbf{B}^D = (B_r, B_\phi, B_z)$ within the disk, having a field strength $\mathbf{B}^D = \mathbf{B}^D(r)$ that depends on the distance r from the neutron star in the disk plane. [10,11] The source that ultimately powers the field B will be the neutron star's gravitational energy.

In a Keplerian disk of finite thickness the accreting matter will have a velocity vector $\mathbf{V} = (-v_r, v_\phi, v_z)$ in a cylindrical coordinate system referred to the symmetry plane of the disk; here v_r is the radial accretion velocity, v_ϕ is the Keplerian velocity $(GM/r)^{1/2}$, and v_z represents the "convective" velocity of flow perpendicular to the disk. The dynamo mechanism will drive processes that build up the field components in the following sequence[12]: $B_r \xrightarrow{v_\varphi} B_\varphi \xrightarrow{v_z} B_z \xrightarrow{v_r} B_r$.

One should recognize that the resultant field outside the disk will remain poloidal in character, with components differing somewhat in strength from the corresponding ones inside the disk. In particular the disk material will constitute hot ($T \sim 10^8$ K) plasma of high conductivity, [13] $\lambda \sim 10^{21}$ cm^{-1}, so that the lines of force of the field $\mathbf{B}^D$ will be frozen into the disk plasma.

2. By Ohm's law, the motions of the plasma and the presence of a magnetic field will serve to generate electric fields E^D in the accretion disk. To estimate the efficiency of this process, we jointly consider Ohm's law and Ampere's law:

$$\begin{cases} \mathbf{I} = \sigma\left[\mathbf{E} + \frac{1}{c}(\mathbf{v}\times\mathbf{B})\right], \\ \mathbf{I} = \frac{c}{4\pi}\,\mathrm{curl}\,\mathbf{B}. \end{cases} \quad (1)$$