Rotation law and magnetic field for M dwarf models

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In stellar convection zones and fully convective stars, the rotation profiles are determined by the balance between the Reynolds stress and the meridional circulation. Due to the Coriolis force, the Reynolds stress has a non-diffusive component called Λ-effect that drives both differential rotation and meridional motions. The solar differential rotation pattern is almost perfectly reproduced by a mixing-length model of the convection zone that takes into account the influence of the Coriolis force on the convective motions. The same model also yields the turbulent electromotive force that together with rotational shear drives the solar dynamo.

The model has recently been applied to a fully convective pre-main sequence star. We find that for a strictly spherical star without any latitudinal gradients in temperature, density and pressure the rotation is very close to the rigid-body state. We conclude that the stellar magnetic field must be generated by a mechanism quite different from that in the Sun, namely an α^2 rather than an $\alpha\Omega$-dynamo. It is thus very likely to have non-axisymmetric geometry and not to show cyclic behavior.

We study the analogous problem for M dwarfs. Like the T Tauri stars, these objects are fully convective and may hence be expected to have similar rotational profiles and magnetic field structures, respectively. As their Coriolis numbers are, however, closer to solar values than to those of pre-main sequence stars, the rotation may also be of solar-type.

1. Introduction

At present, the idea is widely accepted, that the solar magnetic field is generated in the overshoot layer at the bottom of the convection zone. This model has replaced the older one of a convection zone dynamo for three reasons. First, helioseismology has revealed that the radial rotational shear necessary to produce the toroidal field is only found in the overshoot layer. Moreover, the shear is positive and hence an $\alpha\Omega$-type dynamo with the α-effect acting in the convection zone would produce a poleward drift of the magnetic field. Second, in the convection zone the magnetic field is subject to buoyancy forces that transport the flux from the bottom to the surface within about one month and thus inhibit the generation of a large-scale field. A third problem for the convection zone dynamo is the fact that the orientation of the field that emerges from bipolar regions is almost parallel to the equatorial plane, which suggests that the origin of the field is a region that is less turbulent than the convection zone.

While the overshoot dynamo explains the solar cycle quite well, it can obviously not work in fully convective stars. We are thus expecting the M dwarf activity as substantially different from the solar activity. In opposition to the solar case fully convection stars can answer the question whether convection of rotating stars is able to produce an alpha-effect or not. That the answer is Yes in the moment we only know for the case which are simulated by the geophysicists (Glatzmaier & Roberts 1996).

As well-known the internal stellar rotation is one of the key processes for the amplification of magnetic fields. Consequently, we have to know the stellar rotation law in order to design a particular dynamo model. We therefore present here first results for a theory of differential rotation in cool M stars.

2. Theory of differential rotation

In a convective star the turbulent convective motions cause an additional stress on the mean (global) motion known as Reynolds stress. While this stress can be described as an additional viscosity in case of a non-rotating convection zone this is no longer correct as soon as the rotation period becomes comparable with the convective turnover time. In that case a non-viscous contribution, the Λ-effect, arises forcing differential rotation.

In Küker *et al.* (1993), Reynolds stress was considered the only transporter of angular momentum in the solar convection zone. This model yields a rotation pattern that agrees almost perfectly with the observations from helioseismology and thus confirms the assumption, that the meridional flow is of minor importance for the problem of solar differential rotation. The latter assumption can, however, not be correct for stars in general since a model based on Reynolds stress alone always yields a normalized differential rotation that increases with increasing rotation rate, in contradiction to the observations.

We therefore present a model of a rapidly rotating fully convective main sequence star in which Reynolds stress and meridional circulation are treated consistently, i.e. we solve the full Reynolds equation,

$$\rho\left[\frac{\partial\bar{\boldsymbol{u}}}{\partial t}+(\bar{\boldsymbol{u}}\cdot\nabla)\bar{\boldsymbol{u}}\right]=-\nabla\cdot(\rho Q)-\nabla\bar{p}+\rho\boldsymbol{g}+\nabla\cdot\pi. \tag{2.1}$$

Here

$$Q_{ij}=\langle u_i'(\boldsymbol{x},t)u_j'(\boldsymbol{x},t)\rangle \tag{2.2}$$

is the correlation tensor of the fluctuating part $\boldsymbol{u}'$ of the velocity field, $\bar{\boldsymbol{u}}$ denotes its mean velocity. The molecular stress tensor π can be neglected since it is many orders of magnitude smaller than the Reynolds stress.

We only treat the axisymmetric case. The velocity field can then be separated into a global rotation and the meridional flow:

$$\bar{\boldsymbol{u}}=r\sin\theta\,\Omega\,\boldsymbol{e}_\phi+\boldsymbol{u}^{\mathrm{m}}, \tag{2.3}$$

where $\boldsymbol{e}_\phi$ is the unit vector in the azimuthal direction. The azimuthal component of the Reynolds equation reads:

$$\frac{\partial\rho r^2\sin^2\theta\,\Omega}{\partial t}+\nabla\cdot\boldsymbol{t}=0, \tag{2.4}$$

where

$$\boldsymbol{t}=r\sin\theta\left[\rho r\sin\theta\Omega\boldsymbol{u}^{\mathrm{m}}+\rho\langle u_\phi'\boldsymbol{u}'\rangle\right]. \tag{2.5}$$

The meridional circulation can be obtained by taking the azimuthal component of the curl of (2.1):

$$\frac{\partial\omega}{\partial t}=-\left[\nabla\times\frac{1}{\rho}\nabla(\rho Q)\right]_\phi+r\sin\theta\frac{\partial\Omega^2}{\partial z}+\frac{1}{\rho^2}(\nabla\rho\times\nabla p)_\phi, \tag{2.6}$$

where $\omega=(\nabla\times\bar{\boldsymbol{u}})_\phi$. In (2.6), we have omitted all nonlinear terms except the one including Ω^2. We use the anelastic approximation

$$\nabla\cdot(\rho\bar{\boldsymbol{u}})=0, \tag{2.7}$$

i.e. the density is constant with time but varies with depth. Both density and pressure are assumed to be functions of the fractional stellar radius only and their gradients are thus aligned. As a consequence, the last term in (2.6) vanishes and the rotation pattern is determined by the balance between the meridional flow and the Reynolds stress.

In the correlation tensor Q, a viscous and a non-viscous part can be distinguished,

$$Q_{ij} = -\mathcal{N}_{ijkl}\frac{\partial \bar{u}_k}{\partial x_l} + \Lambda_{ijk}\Omega_k. \tag{2.8}$$

The viscous part is given by

$$\begin{aligned}\mathcal{N}_{ijkl} &= \nu_1(\delta_{ik}\delta_{jl} + \delta_{jk}\delta_{il}) \\ &+ \nu_2\left(\delta_{il}\frac{\Omega_j\Omega_k}{\Omega^2} + \delta_{jl}\frac{\Omega_i\Omega_k}{\Omega^2} + \delta_{ik}\frac{\Omega_j\Omega_l}{\Omega^2} + \delta_{jk}\frac{\Omega_i\Omega_l}{\Omega^2} + \delta_{kl}\frac{\Omega_i\Omega_j}{\Omega^2}\right) \\ &+ \nu_3\delta_{ij}\delta_{kl} - \nu_4\delta_{ij}\frac{\Omega_k\Omega_l}{\Omega^2} + \nu_5\frac{\Omega_i\Omega_j\Omega_k\Omega_l}{\Omega^4}\end{aligned} \tag{2.9}$$

(Kitchatinov *et al.* 1994). The viscosity coefficients,

$$\nu_n = \nu_0\phi_n(\Omega^*), \qquad n = 1\ldots 5, \tag{2.10}$$

depend on the angular velocity as well as on the convective turnover time, $\tau_{\rm corr}$, via the Coriolis number

$$\Omega^* = 2\tau_{\rm corr}\Omega. \tag{2.11}$$

In the limiting case of very slow rotation, $\Omega^* \ll 1$, the viscous stress becomes isotropic and reduces to the well-known stress-strain relation with the viscosity coefficient

$$\nu_0 = c_\nu\tau_{\rm corr}u_{\rm T}^2, \tag{2.12}$$

which we therefore use as the reference value of the turbulence viscosity. In (2.12), c_ν is a dimensionless number smaller than unity, $\tau_{\rm corr}$ the correlation time of the turbulence and $u_{\rm T} = \sqrt{\boldsymbol{u}'^2}$ the amplitude of the velocity fluctuations. In stellar convection zones $\tau_{\rm corr}$ equals the convective turnover time while $c_\nu \sim 1/3$.

The second contribution, the Λ-effect, is the source of differential rotation. In spherical polar coordinates, it is only present in the components $Q_{r\phi}$ and $Q_{\theta\phi}$, i.e.

$$Q^\Lambda_{r\phi} = \nu_0(V^{(0)} + V^{(1)}\sin^2\theta)\sin\theta\Omega, \tag{2.13}$$

$$Q^\Lambda_{\theta\phi} = \nu_0 H^{(1)}\sin^2\theta\cos\theta\Omega. \tag{2.14}$$

The functions ϕ_n have been derived in Kitchatinov *et al.* (1994) and those for $V^{(0)}$, $V^{(1)}$, and $H^{(1)}$ can be found in Kitchatinov & Rüdiger (1993). Both the diffusive and non-diffusive terms in the Reynolds stress are suppressed and deformed by rotation. From our point of view the question is how the effects combine to the resulting rotation law.

3. The alpha-puzzle

The evolution of the mean magnetic field $\bar{\boldsymbol{B}}$ is governed by the dynamo equation

$$\frac{\partial\bar{\boldsymbol{B}}}{\partial t} = \nabla\times(\bar{\boldsymbol{u}}\times\bar{\boldsymbol{B}} + \boldsymbol{\mathcal{E}}), \tag{3.15}$$

where $\boldsymbol{\mathcal{E}}$ is the turbulent electromotive force, $\boldsymbol{\mathcal{E}} = \langle\boldsymbol{u}'\times\boldsymbol{B}'\rangle$. As usual, we assume approximate scale separation and write

$$\mathcal{E}_i = \alpha_{ij}\bar{B}_j + \eta_{ijk}\bar{B}_{j,k}. \tag{3.16}$$

Basic for the theory is the knowledge of the tensors α and η. For slow rotation and weak magnetic field these tensors take the simple and well-known forms

$$\alpha_{ij} = \alpha_0\delta_{ij} \quad \text{and} \quad \eta_{ijk} = \eta_{\rm T}\epsilon_{ijk}, \tag{3.17}$$

hence

$$\boldsymbol{\mathcal{E}} = \alpha_0\bar{\boldsymbol{B}} - \eta_{\rm T}\,{\rm rot}\,\bar{\boldsymbol{B}}. \tag{3.18}$$

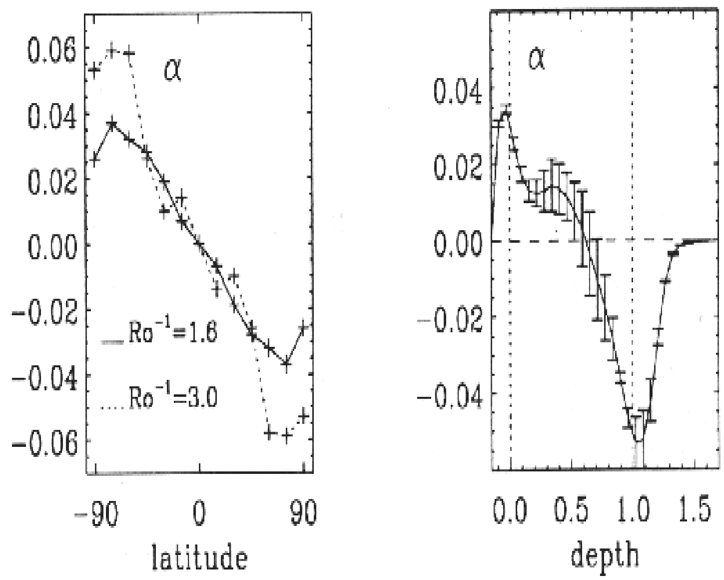

FIGURE 1. LEFT: The latitudinal dependence of α for 2 simulations. RIGHT: The depth dependence of α for the north pole. Note the negative values at the bottom of the convective domain (Brandenburg 1994)

In case of arbitrary rotation rate but weak magnetic field, both the α-effect and the magnetic diffusivity tensor become strongly anisotropic. The α-effect consists of the contributions from the stratifications of density and turbulence velocity. The contribution from density stratification reads

$$\alpha_{ij}^{\rho} = -\delta_{ij}(\boldsymbol{G\Omega})\alpha_1^{\rho} - (G_i\Omega_j + G_j\Omega_i)\alpha_2^{\rho} - (G_i\Omega_j - G_j\Omega_i)\alpha_3^{\rho} - \frac{\Omega_i\Omega_j}{\Omega^2}(\boldsymbol{G\Omega})\alpha_4^{\rho}, \quad (3.19)$$

where $\boldsymbol{G} = \nabla \log \rho$. A similar expression holds for the α-effect from the stratification of the turbulence, i.e.

$$\alpha_{ij}^{u} = -\delta_{ij}(\boldsymbol{U\Omega})\alpha_1^{u} - (U_i\Omega_j + U_j\Omega_i)\alpha_2^{u} - (U_i\Omega_j - U_j\Omega_i)\alpha_3^{u} - \frac{\Omega_i\Omega_j}{\Omega^2}(\boldsymbol{U\Omega})\alpha_4^{u}, \quad (3.20)$$

with $\boldsymbol{U} = \nabla \log \sqrt{u_{\mathrm{T}}^2}$. The vectors $\boldsymbol{G}$ and $\boldsymbol{U}$ point to opposite directions. Hence, the terms (3.19) and (3.20) have different signs. While in the convectively stable overshoot layer at the bottom of the solar convection zone the total sign of α is that of α^u, the α-effect due to the density stratification dominates in stellar convection zones. A widely used approximation for the α-effect is its relation to the helicity $\mathcal{H} = \langle \boldsymbol{u}' \operatorname{rot} \boldsymbol{u}' \rangle$:

$$\alpha \propto -\tau_{\mathrm{corr}}\mathcal{H} \quad (3.21)$$

(Krause & Rädler 1980). Models of SN explosion indeed seem to confirm this relation (Korpi *et al.* 1998). Recently, the helicity has also been derived from solar surface mesogranulation pattern observations as negative (left-handed, see Rüdiger *et al.* 1998). The expected α-value in convection zones is thus positive at the northern hemisphere.

For slow rotation α is proportional to Ω. The numerical simulations are in agreement with this finding. The α-effect results as positive at the top of the convection zone of the northern hemisphere (Fig. 1, Brandenburg 1994). This is in accordance with the

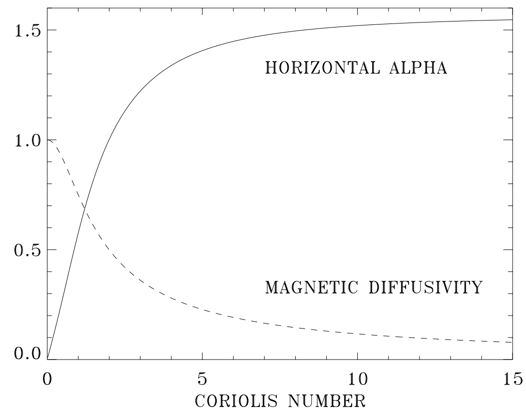

FIGURE 2. The horizontal part of the α-effect, $\alpha_{\phi\phi}$, and the magnetic diffusivity vs. the Coriolis number. Both functions are given in dimensionless units

well-known α-relation

$$\alpha_{\phi\phi} = -l_{\rm corr}^2 \Omega \frac{d}{dr} \log(\rho u_{\rm T}) \cos\theta, \tag{3.22}$$

where the helicity is expressed by the influence of rotation on stratified turbulence. In (3.22), $l_{\rm corr}$ denotes the mixing length. Two remarks are reasonable:

i) The vertical component α_{zz} can have the opposite sign to the horizontal components of the α-tensor (Brandenburg *et al.* 1990). Investigating the expansion flow caused by SN explosions in the interstellar medium Ferrière (1993) also found that α_{zz} can have the opposite sign to the horizontal components of the α-tensor. A negative α_{zz} has also been obtained by Kaisig *et al.* (1993) using an axisymmetric compressible simulation. Rüdiger & Kitchatinov (1993) found a negative value for α_{zz} for stratified turbulence and intermediate values of the inverse Rossby number.

ii) At the base of the convection zone we have a steep *decrease* of the turbulence intensity so that (3.22) turns there to negative values (at northern hemisphere) (Krivodubskij & Schultz 1993). Again, the simulation presented in Fig. 1 is in accord with this argument. Negative values for the α-effect are necessary for the operating of the solar dynamo (Rüdiger & Brandenburg 1995).

There is a new discussion with the α-effect in shear flows. A quasilinear computation (SOCA) of the influence of the differential rotation on the α-effect leads to

$$\alpha \simeq -l_{\rm corr}^2 \frac{\partial\Omega}{\partial\theta} \frac{d\log(\rho u_{\rm T})}{dr} \sin\theta \tag{3.23}$$

for a sphere or

$$\alpha \simeq -l_{\rm corr}^2 \frac{d\Omega}{ds} \frac{d\log\rho}{dz} \tag{3.24}$$

for an accretion disk (Pipin *et al.* 1999), where s is the distance from the axis of rotation. The latter relation for accretion disks with $\partial\Omega/\partial s < 0$ yields negative values in the northern hemisphere and positive values in the southern hemisphere (Brandenburg & Donner 1997). For the solar overshoot region, however, with (3.23) negative values for the northern α only result for the steep decrease of the turbulence intensity with depth. The majority of the turbulence quantities in the mean-field theory vanishes for high rotation rates (Kitchatinov *et al.* 1994). Interestingly enough it is not true for the

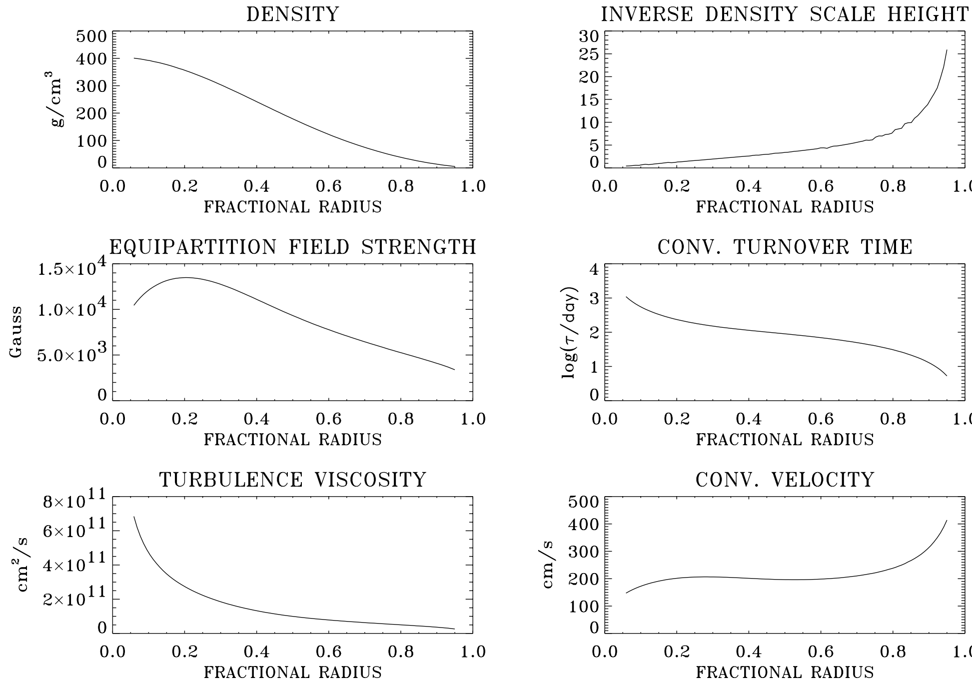

FIGURE 3. The stratification of an M dwarf with $0.1M_\odot$. All quantities are plotted vs. the fractional stellar radius.

α-effect itself. For fast rotation the α-effect saturates but becomes highly anisotropic (Rüdiger 1978; Rüdiger & Kitchatinov 1993).

In Fig. 2, the dependences of the horizontal part of the α-effect, $\alpha_{\phi\phi}$, and the magnetic diffusivity on the Coriolis number are shown. In case of rapid rotation, i.e. large Coriolis numbers, $\alpha_{\phi\phi}$ approaches a finite limit while the magnetic diffusivity decreases as $1/\Omega^*$. For sufficiently rapid rotation the dynamo will thus always become supercritical.

4. Models and their internal flows

Stellar structure models by Chabrier and Baraffe (1997) are used, where all stars with less than 0.35 solar masses were found to be fully convective. The first model describes a star with 0.3 solar masses at an age of 5.6 Gyr. This star has a radius of 0.3 solar radii and a central density of 90 g/cm^3. The second star has 0.1 solar masses, an age of 6.6 Gyr, 0.12 solar radii, and a central density of 400 g/cm^3. Fig. 3 shows its stratification. From top left to bottom right, the plots show density, amplitude of the density stratification vector, $\boldsymbol{G} = \log \nabla\rho$, equipartition field strength $B_{\rm eq}$, convective turnover time $\tau_{\rm corr}$, reference value ν_0 of the turbulence viscosity and convection velocity $u_{\rm T}$ from mixing-length theory. While the radius of this star approximately equals the depth of the convection zone of the Sun, the densities are two orders of magnitude larger. The convection velocities, on the other hand, are much smaller than those in the Sun. The convective turnover times of about three months result in Coriolis numbers larger than one, especially for young stars with rotation periods of several days or even less than one day. Hence, in the context of mean-field magnetohydrodynamics, M dwarfs must be regarded as rapid rotators. Slow rotation is only reached for M dwarfs with rotation periods exceeding 1

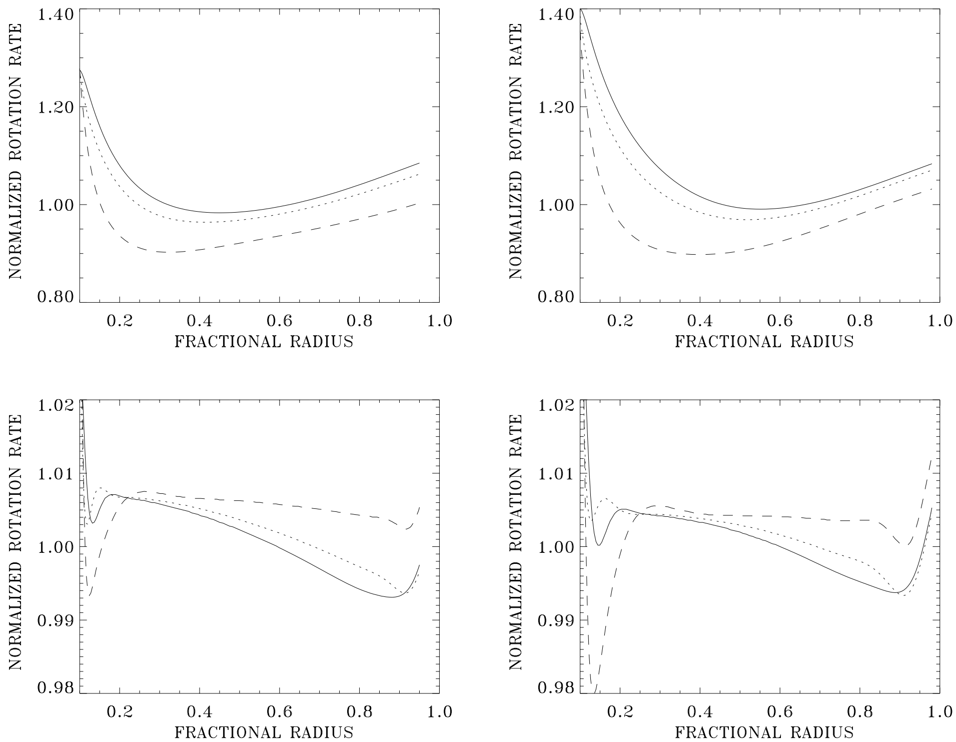

FIGURE 4. Normalized rotation rates at the equator (solid lines), 30° latitude (dotted), and 60° latitude (dashed) for M dwarfs with $M = 0.1 M_\odot$ (left column) and $M = 0.3 M_\odot$ (right column). The average angular velocities are $10^{-7}\mathrm{s}^{-1}$ (top) and $10^{-6}\mathrm{s}^{-1}$ (bottom).

year.† To determine the rotation profiles, we have solved the system (2.4,2.6) for the stellar models described above with an explicit time dependent finite difference code in spherical polar coordinates. The stellar surface is assumed to be stress-free,

$$Q_{r\phi} = Q_{r\theta} = 0. \tag{4.25}$$

For technical reasons, we exclude the innermost part of the star by imposing a second boundary at a fractional radius of 0.1 with the same boundary conditions.

Fig. 4 shows the resulting rotation profiles for two different angular frequencies, namely $\Omega = 10^{-7}\mathrm{s}^{-1}$ and $\Omega = 10^{-6}\mathrm{s}^{-1}$. The first value corresponds to very slow rotation with a period of 727 days while in the second case the period of 72 days is still three times longer than that of the solar rotation. In the first case (top row), the Coriolis numbers are larger than one at small radii and smaller than one close to the surface. Hence, there are different rotation patterns. While the central parts of the stars show negative radial shear, the rotation profiles are flatter and the shear is positive in the upper layers. The plots in the bottom row show a completely different picture. The rotation has become almost rigid. There is now a very small negative shear throughout the whole star except two thin layers at the boundaries and the interior of the tangent cylinder around the (artificial) inner boundary. For more rapid rotation the stars continue to rotate more and more rigidly.

† Proxima Cen: 31.5 days, Lalande 21185: 47 days

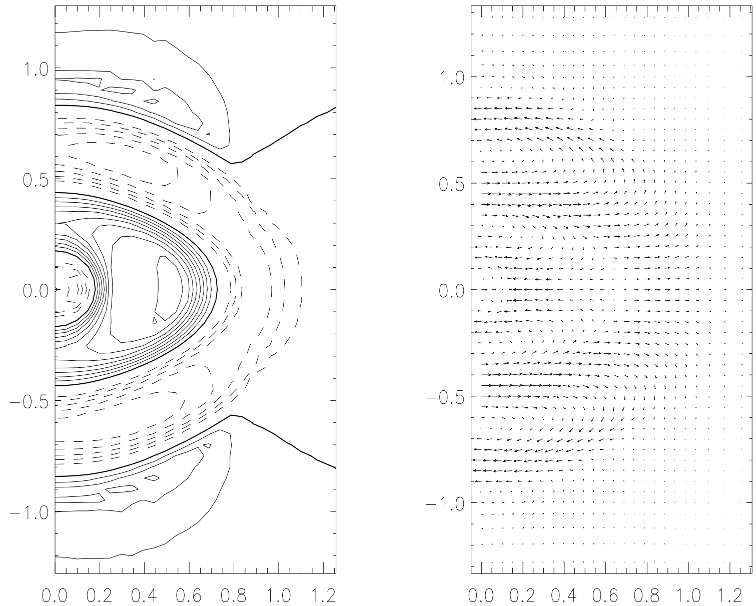

FIGURE 5. Magnetic field structure of an M dwarf with 0.1 $M_\odot$. The left picture shows the azimuthal component of the field in the $\phi = 0$ half plane. Solid contours denote positive and dashed contours negative values of B_ϕ. The right picture is a vector plot of the field components in the same half plane. The stellar radius has been used as unit length.

5. The M star dynamo model

After all, the dynamo operating in the fully convective M dwarfs should clearly be of the α^2-type. This seems as a simplification but it is not. In α^2-dynamos the anisotropy of the α-tensor plays an important role (Rüdiger & Elstner 1994). If – as it is the result of quasilinear turbulence theory – the α_{zz} has the opposite sign as the $\alpha_{\phi\phi}$ – the resulting magnetic field becomes nonaxisymmetric with massive consequences for nonlinear computations (cf. Moss & Brandenburg 1995). The anisotropy of the α-effect rather than its sign gives the main uncertainty in the theory. The growth of the mean magnetic field is limited by its back reaction on the electromotive force. We use a simple α-quenching prescription,

$$\alpha_0 = \frac{\tilde{\alpha}_0}{1+\beta^2}, \tag{5.26}$$

where

$$\beta = 2\pi\,\frac{|B|}{B_{\rm eq}} \tag{5.27}$$

is the normalized magnetic field and

$$B_{\rm eq} = \sqrt{\mu_0 \rho u_{\rm T}^2}. \tag{5.28}$$

is the turbulence-equipartition field.

The magnetic diffusivity tensor assumes the form

$$\eta_{ijk} = \eta_{\rm T}\epsilon_{ijk} + (\eta_\| - \eta_{\rm T})\epsilon_{ijl}\frac{\Omega_l\Omega_k}{\Omega^2} \tag{5.29}$$

with

$$\eta_{\rm T} = (\phi_1 + \phi_2)\eta_0, \qquad \eta_{\|} = 2\phi_1\eta_0 \tag{5.30}$$

and

$$\phi_1 = \frac{3}{4\Omega^{*2}}\left(-1 + \frac{\Omega^{*2}+1}{\Omega^*}\arctan\Omega^*\right), \tag{5.31}$$

$$\phi_2 = \frac{3}{2\Omega^{*2}}\left(1 - \frac{\arctan\Omega^*}{\Omega^*}\right) \tag{5.32}$$

(Kitchatiov *et al.* 1994). The reference value η_0 of the eddy diffusivity is

$$\eta_0 = c_\eta u_{\rm T}^2 \tau_{\rm corr} \tag{5.33}$$

with $c_\eta \simeq 0.30$. To solve the induction equation, we use a 3D time dependent fully explicit second order finite-difference scheme in cylindrical polar coordinates. The computational domain is a cylinder with a radius of two and a height of four stellar radii. The star is located at the center of the cylinder. The surrounding medium is considered co-rotating with the star and poorly conducting. Comparisons with Moss & Brandenburg (1995) show that with the same electromotive force, our model yields very similar results to their model, which assumed a star surrounded by vacuum. The assumption of small conductivity (large magnetic diffusivity) outside the star is thus a rather perfect approximation for a vacuum boundary condition. As there is no rotational shear in our model, the stellar rotation rate enters only via the Coriolis number which we thus vary as an input parameter. There are no further free parameters. For both stellar models the critical value for dynamo excitation lies between one and two. The field geometry for a fairly supercritical case with $\Omega^* = 2$ is shown in Figs. 5 and 6. The field is completely non-axisymmetric without any axisymmetric contribution and symmetric with respect to the equatorial plane. The total field energy is constant with time but the field rotates with the star. The cross sections with constant z show a spiral-type geometry while the field distribution in the r-z half plane is rather shellular. For larger values of the Coriolis number the field resembles that of a tilted dipole with the dipole axis lying in the equatorial plane. This type of field geometry was also found by Moss & Brandenburg (1995). The similarity of the results is due to the fact that α^2-dynamos are not sensitive to the sign of α.

6. Conclusions

Very slow rotators and very fast rotators are almost rigid rotators. As M dwarfs with rotation periods shorter than 1 year must be considered as very fast rotators, their differential rotation proves to be small. There is thus no doubt that the dynamo operating in the fully convective M dwarfs is of the α^2-type. The anisotropy of the α-tensor, suggested by both quasilinear as well as nonlinear simulations, finally leads to nonaxisymmetric and non-oscillating large-scale magnetic fields (cf. Fig. 7). Hence, a distinct rotational modulation should be observable for the active M dwarfs. Our scenario, supported by nonlinear simulations of internal rotation, meridional flow and mean magnetic fields leads to clear predictions for observations. Only very slow rotators with $P_{\rm rot} > 1$ year should adopt solar-type solutions. It is an open question, however, whether the magnetic spin-down of M dwarfs is effective enough to produce such rotation times within the age of the universe.

Our warmest thanks go to Isabelle Baraffe and Gilles Chabrier for their kind cooperation concerning structure and evolution of the used M dwarf models

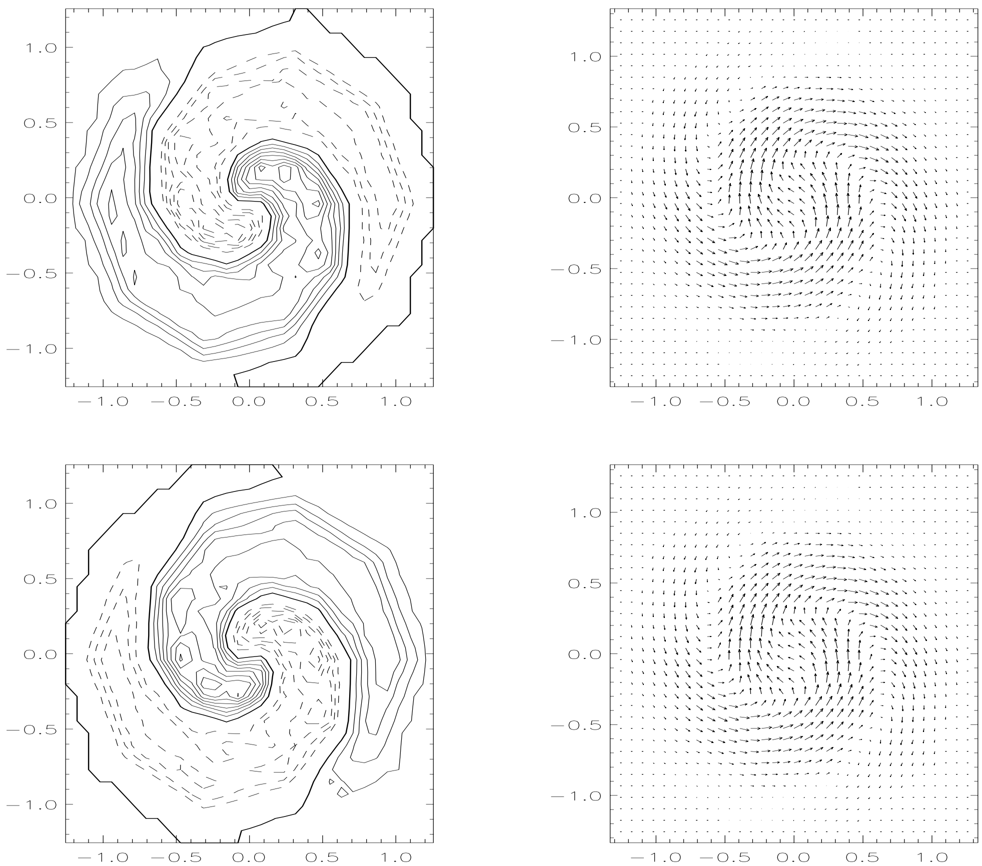

FIGURE 6. The magnetic field of an M dwarf with 0.1 $M_\odot$ in two planes parallel to the equatorial plane. The pictures in the top row show a plane with $z > 0$, while the pictures in the bottom row show a plane with $z < 0$. The pictures in the left column show contour plots of the z component, those in the right column the r and ϕ components.

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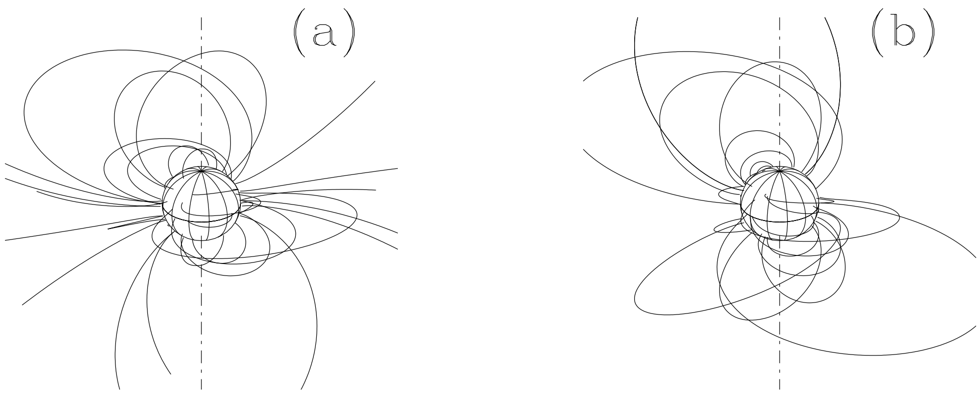

FIGURE 7. Visualization of a magnetic field configuration similar to the one we find by Moss & Brandenburg (1995).

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