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ABSTRACT It is by now well known that most main sequence stars continuously lose mass as a consequence of the winds they emit. In addition to affecting the thermal and dynamical state of the stellar atmosphere, such mass loss can also induce changes in the interiors of stars. In the present review, we consider a few of the ways in which sustained, wind-like mass loss can alter the physical state of main sequence stellar interiors by examining the differences in internal structure, composition, and rotation between mass-losing and conservatively evolving stars.

INTRODUCTION

Numerous observational and theoretical results obtained over the past quarter century have contributed to the current realization that mass loss is a relatively common phenomenon, occurring among stars located nearly everywhere in the H-R diagram. There now exists an abundance of evidence (some of it direct, some of it circumstantial) suggesting that virtually all stars, regardless of effective temperature or evolutionary status, sustain mass loss in the form of a wind. As a research subject mass loss is frequently studied with the goal of identifying and/or elucidating the physical processes responsible for the initiation, acceleration, and maintenance of observed stellar outflows. However, in keeping with the theme of the present meeting, we herein consider how winds can affect the *interiors* of stars, specifically, stars on or near the main sequence.

In an effort to depict a variety of possible effects yet keep the analysis simple, we restrict our attention to three representative examples of the consequences of mass loss for main sequence (MS) stellar interiors. In the first, we use the evolution of massive stars ($M_* \gtrsim 15 M_\odot$) during core hydrogen burning to illustrate how the removal of a significant fraction of the initial stellar mass by a wind can induce internal, structural changes. In the second, we examine the advective effect of mass loss by considering how winds can moderate surface abundance anomalies produced by the diffusive separation of

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chemical elements in stars of intermediate mass ($1.5M_\odot \lesssim M_* \lesssim 3.0M_\odot$). Finally, in the third example, we describe how the torque associated with mass loss from a rotating star can affect the physical state of the interior by considering how the internal solar rotation accommodates the shear which is continuously introduced by the wind-related braking of the sun's surface layers.

MS EVOLUTION OF MASSIVE STARS WITH MASS LOSS

Massive, hot stars having spectral types O through early B and belonging to all luminosity classes are observed to lose mass at rates in the range $10^{-8} \lesssim \dot{M}_* \lesssim 10^{-5} M_\odot\ yr^{-1}$ (see e.g., Cassinelli and Lamers 1987). Such mass loss is thought to be driven by the force produced when radiation from the stellar photosphere is absorbed or scattered in UV spectral lines formed in the flow. The inferred $\dot{M}_*$ values are of sufficient magnitude that the internal structure of an early-type star undergoing mass loss may change over the course of its MS lifetime τ_{MS}, as a consequence of the changing stellar mass. For example, a star with $M_* = 30M_\odot$ which loses mass at the constant rate $\dot{M}_* = 5 \times 10^{-7} M_\odot$ yr^{-1} will suffer a fractional reduction $\dot{M}_*\ \tau_{MS}/M_* = 0.1$ of its initial mass, for $\tau_{MS} = 6 \times 10^6$ years. The structural and evolutionary modifications to constant mass models arising from MS mass loss of this kind have been examined by numerous authors. Among the investigations most pertinent to the present discussion are those by Tanaka (1966), Hartwick (1967), Chiosi and Nasi (1974), Dearborn *et al.* (1978), Falk and Mitalas (1981), Brunish and Truran (1982), and Mitalas and Falk (1984). Additional references are contained in the reviews by Maeder (1984) and Chiosi and Maeder (1986).

To discern the ways in which the physical state of the interior of a mass-losing star differs from that of a conservatively evolving star, we utilize the simple, semi-analytic model described by Falk and Mitalas (1981, hereafter FM; see also Mitalas and Falk 1984; Stein 1966). From comparisons with the results of detailed model calculations, these authors found that during the phase of core hydrogen burning, the relation between stellar mass M_* and luminosity L_* could be accurately represented by

$$(L_*/L_\odot) = A\ \beta^4\ <\mu>^4\ (M_*/M_\odot)^3\ . \tag{1}$$

In equation (1), β is the ratio of gas to total pressure, $<\mu>$ is the mean molecular weight averaged over the mass of the star, and A is a constant whose value depends upon the initial stellar mass and composition.

FM further noted that the contraction of the convective core that occurs as MS evolution proceeds leads to an internal composition profile with the following characteristics. At a given time $t(<\tau_{MS})$, the hydrogen mass fraction X in the outer envelope (i.e., the region external to the initial core boundary) has the uniform value X_0. In the core, X is also uniform, with value $X_c(< X_0)$, reflecting the nuclear processing which has taken place prior to the time in question. Throughout the region comprised of material which was previously a part of the core, X has very nearly a linear dependence on mass. Following FM, we assume that the composition profile within this intermediate region is strictly linear, making possible direct evaluation of the average $<\mu(X)>$ in equation (1). The quantity β can then be obtained by solving the Eddington quartic

$$1-\beta = a\,\beta^4\ <\mu>^4\ (M_*/M_\odot)^2\ , \tag{2}$$

where $a = 0.00309$ (cf. FM).

The fractional mass contained within the convective core can be estimated by assuming that the entire stellar luminosity is produced therein, and utilizing the fact that radiative and adiabatic gradients are equal at the core boundary (cf. Stein 1966). If electron scattering constitutes the primary source of opacity, it can be shown that

$$q \equiv \frac{M_{core}}{M_*} = q_0\ \left(\frac{1+X_c}{1+X_0}\right)\ , \tag{3}$$

where the initial core mass q_0 is a known quantity, expressible in terms of specified constants characterizing the model (cf. FM). With this result, the time-evolution of X_c can be determined by integrating

$$\frac{dX_c}{dt} = -\frac{L_*}{EM_{core}}\ \propto\ \left(\frac{1-\beta}{1+X_c}\right)\ , \tag{4}$$

where $E(\approx 6\times10^{18}$ erg g$^{-1})$ is the energy released from 1 gram of material in which hydrogen has been transmuted to helium via the CNO cycle.

The final component of the model is an equation governing the time evolution of the stellar mass. In the present paper, we consider four distinct prescriptions for the mass loss rate, enumerated below:

$$
\begin{aligned}
dM_*/dt &= 0 && (5a)\\
&= 1.35\times10^{-7}\ (L_*/10^5 L_\odot)^{1.62}, && (5b)\\
&= 3.26\times10^{-7}\left(\frac{M_*}{M_\odot}\right)\ \frac{\Gamma^{1.49}}{(1-\Gamma)^{0.49}}, && (5c)\\
&= 1.72\times10^{-6}\ \left(\frac{L_*}{10^5 L_\odot}\right)\ \left(\frac{R_*}{R_\odot}\right)\ \left(\frac{M_\odot}{M_*}\right), && (5d)
\end{aligned}
$$

where in 5(c), $\Gamma\ =\ 2.40\ (L_*/10^5\ L_\odot)(M_\odot/M_*)$. The mass loss rates given in equations 5(b) - (d) have units $M_\odot$ yr^{-1}, and are, respectively, an empirical fit to observed values obtained by Garmany and Conti (1984), the theoretical rate derived from the line-driven wind model of Castor, Abbott, and Klein (1975), and a proportionality suggested by Mc Crea (1962), scaled to yield $\dot{M}_*\ = 5\times10^{-7} M_\odot$ yr^{-1} for parameters appropriate to a zero-age MS star with $M_* = 30M_\odot$.

We have traced the MS evolution of a $30M_\odot$ star, including mass loss, by integrating equations (4) and (5), assuming $X_0 = 0.71$. Equations (1) (with $A = 71.278$; cf. FM) and (2) were used to determine β and L_* with $<\mu>$ calculated as described above, assuming $\mu(X)\ =\ 4M_H/(5X+3)$. A summary of results is depicted in Figure 1; in each panel, the lettered curves represent models computed with the corresponding mass loss rate given in equations 5(a)-(d). For each model, the results portrayed in the figure span the time interval over which X_c is reduced from its initial value X_0 to the

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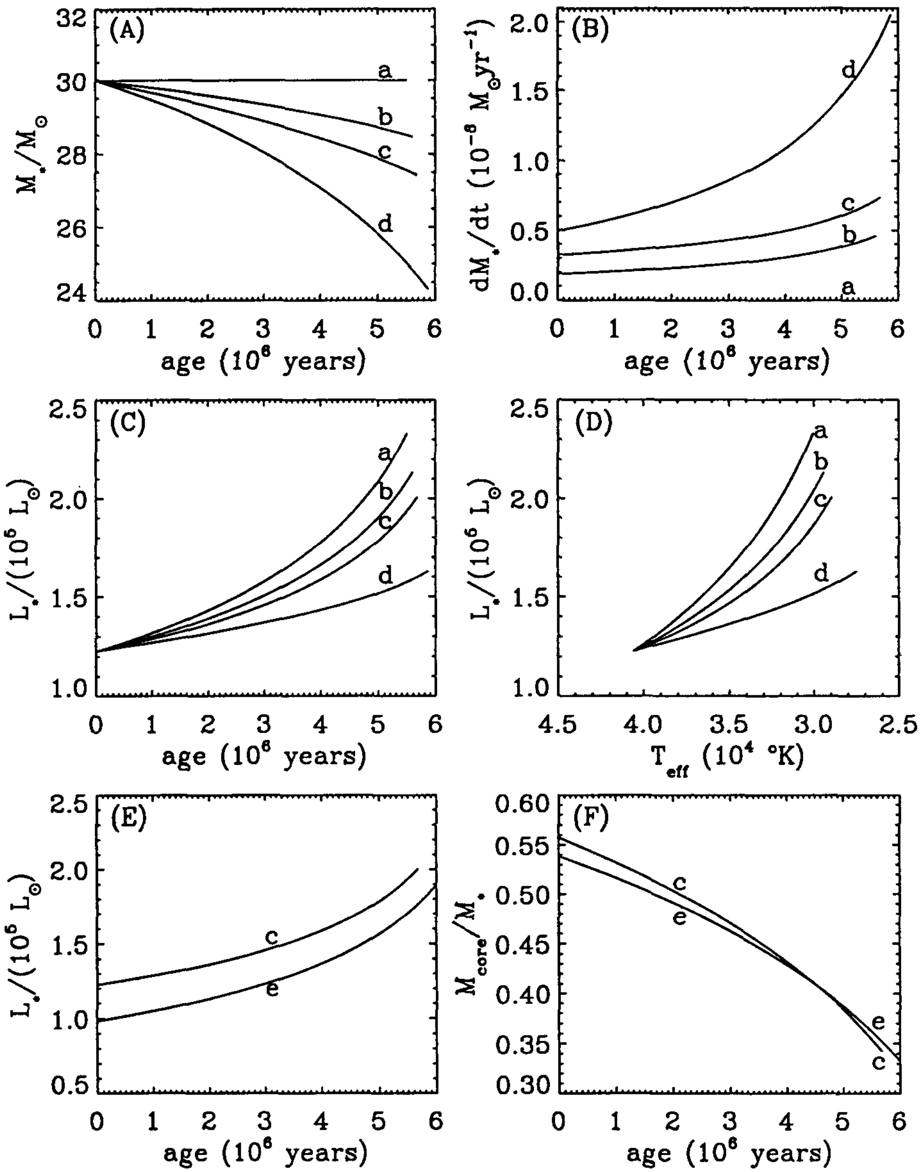

Figure 1: MS evolution of a 30 $M_{\odot}$ star with mass loss, as described in the text.

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value $X_c = 0.05$. In panels (A) - (C), the times required for models (a) - (d) to attain this state of near hydrogen exhaustion in the core are, respectively, $5.50 \times 10^6, 5.60 \times 10^6, 5.68 \times 10^6$, and 5.87×10^6 years. We therefore conclude that the MS lifetime (i.e., the time during which hydrogen is burned in the core) of a mass-losing star is longer than that of a conservatively evolving star. Moreover, as is evident from panels (A) and (B), the rates of mass loss given by equations 5(b), (c), and (d) are (in that order) an increasing sequence at all times. Hence, for the moderate mass loss rates considered here, the enhancement of τ_{MS} (relative to a conservative star of the same initial mass) is larger for larger values of $\dot{M}_*$.

In panel (C), we display the time evolution of the stellar luminosity for each of the models. Note that at a fixed time, the luminosity of a star undergoing mass loss is smaller than that of a star having the same mass at $t = 0$ with $\dot{M}_* = 0$. Because $L_* \propto M_*^3$ (cf. eq. [1]), the model which experiences the most mass loss during MS evolution is characterized by the smallest luminosity at later times. However, since at a given time that model also has the largest value of X_c, it has a commensurably larger core mass fraction q (cf. eq. [3]). At $t = 5 \times 10^6$ years, model (a) has $X_c = 0.150$ and $q = 0.375$, while model (d) has $X_c = 0.197$ and $q = 0.390$. In panel (D), the computed evolutionary tracks are depicted in an H-R diagram. Apparent in the figure is the tendency for the tracks of stars suffering more cumulative mass loss to assume an increasingly horizontal orientation. Note also that as the tracks become less inclined, they extend to cooler effective temperatures. In model (a), $T_{eff} = 30{,}100$ K when $X_c = 0.05$, while in model (d), $T_{eff} = 27{,}500$ K for the same value of X_c. Hence, the MS width in the H-R diagram is larger for a mass-losing star than for a conservative star, other things being equal.

For model (c), $M_* = 27.43 M_\odot$ when $X_c = 0.05$ near the termination of core hydrogen burning. It is of interest to compare the properties of this model with those of a model having $M_* = 27.43 M_\odot$ and $\dot{M}_* = 0$, evolved to the same central hydrogen abundance. This is done in panels (E) and (F), wherein the curves pertaining to the latter, conservatively evolving star are labelled (e). As noted above, model (c) attains $X_c = 0.05$ at $t = 5.68 \times 10^6$ years; in model (e) hydrogen is depleted to this level at $t = 6.00 \times 10^6$ years. Inspection of panel (E) reveals that the luminosity of the mass-losing model exceeds that of the $\dot{M}_* = 0$ model at all times. In particular, L_* for model (c) is larger than L_* for model (e) even when both stars have the same values of M_* and X_c. The primary reason for this behavior can be ascertained from panel (F), which shows the time evolution of the convective core mass fraction for both models. As can be determined from the figure, when $M_* = 27.43 M_\odot$ and $X_c = 0.05, q = 0.342$ in model (c), while $q = 0.331$ for model (e). Thus, in comparison to a star of the same mass and central hydrogen abundance, a star with $\dot{M}_* \neq 0$ has a more massive convective core. Because $L_* \approx \bar{\epsilon} M_{core}$, where $\bar{\epsilon}$ is the mass-averaged specific energy generation rate, the luminosity of model (c) is greater than that of model (e).

We close this section by noting that an additional, well-studied consequence of mass loss on the internal structure of massive stars is its effect on the occurrence, structure, and lifetime of semi-convective regions. The reader is referred to the previously cited papers for discussion of this topic.

ABUNDANCE ANOMALIES IN MS STARS OF INTERMEDIATE MASS

Definitive observational evidence for the existence of winds from MS stars with spectral types A-F is presently unavailable. However, theoretical attempts to account for the chemical peculiarities exhibited by some such stars may provide an indication of the occurrence of mass loss from them.

The stars that constitute the class of metallic-line objects designated "Fm-Am stars" have the following attributes in common (Cayrel *et al.*, 1991; Boyarchuk and Savanov 1986; Preston 1974). They are MS stars with spectral types between about F0 and A0, effective temperatures between 7000 K and 10,000 K, and masses in the range $1.5 \lesssim (M_*/M_\odot) \lesssim 3.0$. Spectroscopic analyses have revealed that approximately 20% of all stars with these properties exhibit anomalous photospheric abundances. The classical pattern of abundance anomalies observed to be characteristic of the Fm-Am stars as a group consists of underabundances (relative to solar photospheric abundances) of the elements Ca and Sc, and overabundances of virtually all elements heavier than Fe. Underabundant species are typically deficient by a factor $\sim$ 5-10, while overabundant species can be enhanced by as much as a factor $\sim$ 10-100, although substantial dispersion exists among both elements and stars. The Fm-Am stars do not have detectable surface magnetic fields, and, when compared with chemically normal stars in the same spectral type range, they have significantly lower rotational velocities. This latter characteristic may be related to the fact that the vast preponderance of metallic-line stars are members of binary systems (Abt and Levy 1985).

The chemical peculiarities observed in Fm-Am stars are thought to be the result of selective, radiatively-driven diffusion of some species in the outer layers of the stellar interior (Michaud 1970; Watson 1970, 1971; see also the reviews by Michaud 1977, 1987, 1991; Vauclair and Vauclair 1982). Stars having these spectral types arrive on the MS in a uniformly mixed state, with subphotospheric convection zones due to both H and He. These adjacent, superficial convective regions are quite thin; for $M_* \geq 1.5 M_\odot$, the mass fraction contained in the outer, H convection zone is $M_{conv}/M_* \lesssim 10^{-9}$ (see, e.g., Michaud *et al.*, 1976). Over time He settles gravitationally; for stars with rotational velocities low enough that meridional circulation is unimportant, after $\sim 10^6 - 10^7$ years of MS evolution, the initial He abundance in the outer envelope is diminished by a factor of about 3, and the associated inner convection zone disappears. Upward diffusion of heavy elements can subsequently take place in the stable layers just beneath the remaining H convection zone.

At the base of the H convection zone, the prevailing physical conditions are such that the dominant ionization stages of many chemical elements have bound-bound transitions with frequencies situated near the peak of the stellar flux distribution. For abundant species (roughly, those lighter than Fe), the upward acceleration arising from the absorption of radiation in such lines is generally smaller than gravity, being limited in magnitude by large line optical depths. Ions of these elements undergo gravitational settling. However, for species whose intrinsic abundances are sufficiently small that they can be accurately described as trace constituents of the gas (i.e., elements heavier than Fe), the radiative acceleration associated with optically thin lines can exceed gravity by a large amount.

Consider an ion of element A with mass $m_A = Am_H$ and concentration $c_A = N_A/N_H \ll 1$. In the absence of large concentration gradients and

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neglecting thermal diffusion, the velocity of diffusion of A relative to the background electron-proton gas is approximately (Chapman and Cowling 1970; Aller and Chapman 1960)

$$u_D \approx -D(g - g_{rad})\ \frac{m_A}{kT} \approx -(g - g_{rad})\ \tau_{coll}\ . \tag{6}$$

In equation (6), $D(\sim\ v^2_{th,A}\ \tau_{coll})$ is the coefficient of diffusion, g and g_{rad} are, respectively, the gravitational and radiative accelerations, τ_{coll} is the A-proton collision time, and $u_D \gtrless 0$ implies motion in the $\pm\hat{e}_r$ direction. According to (6), species for which $g_{rad}\ >\ g$ in the region beneath the H convection zone are pushed upward, while those with $g_{rad} < g$ sink under the influence of gravity. Upwardly diffusing species, upon entering the convection zone, are quickly mixed throughout the volume by the turbulent motions therein. The time evolution of the relative abundance of such an element A in the surface layers of the star is governed by the equation of continuity

$$\frac{\partial}{\partial t}(c_A N_H) + \nabla \cdot (c_A N_H \mathbf{u}_D) = 0. \tag{7}$$

By integrating equation (7) over the volume of the convection zone and applying Gauss theorem, it follows that

$$M_{conv}\ \frac{\partial c_A}{\partial t} = 4\pi\ R^2_{conv}\ (\rho c_A u_D)\Big|_{r=R_{conv}}\ , \tag{8}$$

where ρ is the total mass density, and M_{conv} and R_{conv} are the mass and base radius of the convection zone. Hence, the concentration of element A within the convection zone grows at a rate which depends upon the magnitude of the current diffusive flux into it from below.

The theory outlined above is remarkably successful in reproducing the pattern of abundance anomalies observed in Fm-Am stars. That is, species for which the *computed* g_{rad} is greater than g below the H convection zone are generally *observed* to be overabundant, and vice versa. However, the magnitudes of the abundance enhancements obtained from equations (6) and (8) together with calculated g_{rad} values are often significantly larger than those observed; overabundances by factors in excess of 10^4 can be produced in times $\lesssim \tau_{MS}$ (Michaud *et al.*, 1976). According to (8), the H convection zone functions as a reservoir which, in the absence of additional effects, upward diffusion from deeper layers continuously fills. The rate of filling is decreased only if g_{rad} becomes $< g$ (e.g., by saturation), or if the region in which $g_{rad} > g$ for a particular element becomes depleted of that element.

An alternative possibility for limiting the magnitudes of diffusively produced overabundances is to drain the reservoir from above by mass loss. Assume that a typical star of this type undergoes steady, spherical mass loss at a rate $\dot{M}_*$ which is insufficient to alter the internal stellar structure. Conservation of mass then requires that a fluid parcel at radius r within the star have the small, outward velocity $u_W = \dot{M}_*/(4\pi\ r^2\rho)$. Hence, the continuity equation (7) for species A must be amended by making the substitution $\mathbf{u}_D \rightarrow \mathbf{u}_D + \mathbf{u}_W$, so that integration over the convection zone volume leads to the result

$$M_{conv}\ \frac{\partial c_A}{\partial t} = 4\pi R^2_{conv}\ \ [\rho c_A(u_D + u_W)]\Big|_{r=R_{conv}} - \dot{M}_* c_A\ , \tag{9}$$

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where it has been assumed that the composition of the wind is the same as that of the convection zone. Note that equation (9) contains the sink term $-\dot{M}_* c_A$, describing the rate at which ions of species A are removed from the convection zone as a result of advection by the wind flow. When the evolution of surface abundances for Fm-Am stars is computed using equation (9) instead of (8), it is found that the concentrations of elements having $g_{rad} > g$ beneath the H convection zone grow only for a time $\sim M_{conv}/\dot{M}_*$ before declining at later times. Consider the results of Michaud *et al.* (1983) for a star with $M_* = 1.8M_\odot, T_{eff} = 7800$ K, $\log g = 4.15$, and $M_{conv}/M_* \approx 10^{-9}$. These authors find that for mass loss rates in the range $\dot{M}_* = 10^{-14}-10^{-15} M_\odot$ yr^{-1}, the surface abundances of species such as Ba, Eu, and Hg (for which $g_{rad} > g$) cease to increase some $10^5 - 10^6$ years after the disappearance of the He convection zone. At this stage, the diffusive overabundances of these elements are ~10-50, in accord with observational determinations. Moreover, for such modest rates of mass loss, the wind-induced, upward advection of *all* species (arising from the first term on the RHS of equation (9)) is sufficiently slow that elements having $g_{rad} < g$ below the H convection zone (e.g., Ca and Sc) remain underabundant in the stellar surface layers for times $\sim \tau_{MS}$. Hence, models based on selective, radiative diffusion in the presence of steady mass loss appear to be capable of accounting for both the morphology and magnitude of the observed chemical peculiarities in Fm-Am stars.

As noted at the outset of this section, winds from MS stars in the spectral type range A-F have yet to be directly observed. If the rates of mass loss from these stars are typically as low as the values inferred from the application of diffusion theory to Fm-Am stars, then the current dearth of wind detections is not surprising. In this regard, recent observations at optical and radio wavelengths (Lanz and Catala 1992; Brown *et al.*, 1990) have established upper limits of between 10^{-9} and $10^{-10} M_\odot$ yr^{-1} for the rates of mass loss from a number of A and F dwarf stars.

INTERNAL ROTATION OF SOLAR-TYPE STARS

Direct observations of the solar wind from satellites in Earth orbit have established that the Sun loses mass at a rate $\dot{M}_\odot \sim 10^{12}$ g s^{-1} $\sim 10^{-14} M_\odot$ yr^{-1}. If the current value of $\dot{M}_\odot$ is typical of its magnitude over the entire course of the Sun's MS lifetime, then an amount of mass $\sim 10^{-4} M_\odot$ would be lost during that period, assuming $\tau_{MS} \sim 10^{10}$ years. Such a reduction of the solar mass is too small to engender structural changes of the type discussed previously. Moreover, although a hydrogen convection zone exists in the outer layers of the solar interior, the physical conditions which prevail at its base ($M_{conv}/M_\odot \approx 1.5 \times 10^{-2}, T \approx 2 \times 10^6$ K) are not conducive to the segregation of chemical elements by diffusion, as in the Fm-Am stars (e.g., $M_{conv}/M_* \sim 10^{-9}, T \approx 3 \times 10^4$ K for $M_* = 1.8M_\odot$). There remains, however, one physical property of the Sun's interior which is profoundly affected by the solar wind, namely, its rotational state.

Because the Sun rotates, the wind that emanates from the outer, coronal layers of the solar atmosphere removes angular momentum as well as mass. Near the base of the flow, the solar magnetic field enforces a state of approximate corotation of the plasma with the Sun. The resulting torque exerted on the Sun by the magnetized wind opposes the solar rotation, and is of sufficient magnitude to significantly reduce the solar angular velocity in a time $\lesssim \tau_{MS}$ (Weber and Davis 1967). Spacecraft observations of the wind plasma and interplanetary magnetic field yield estimates in the range $0.2 - 0.3 \times 10^{30}$ dyne cm sr^{-1} for the total angular momentum flux associated with outflow in the solar equatorial plane (Pizzo *et al.*, 1983). Integrating this result over all solid angles, it follows that the rate at which the Sun's rotation is braked by the solar wind is about $\dot{J}_\odot \approx 3 \times 10^{30}$ dyne cm, implying an e-folding time for angular momentum loss (assuming rigid rotation) $J_\odot/\dot{J}_\odot \sim 10^{10}$ years. Note that the application of this torque to the surface layers of the Sun gives rise to stresses within the sheared fluid which act to redistribute the angular momentum contained therein. In the convection zone, the viscosity associated with turbulent fluid motions is sufficiently great (e.g., $\nu \sim 10^{12}$ cm^2 s^{-1}) that the deceleration of the topmost layers is communicated to the bottom in a time $(R_\odot - R_{conv})^2/\nu \sim$ a few years. Hence, angular momentum redistribution in response to rotational braking takes place virtually instantaneously (i.e., in comparison to the angular momentum loss time scale) within the convection zone.

In contrast, the time required for the radiative interior of the Sun to respond to a shear imposed at its outer boundary is much less certain. This is due in part to the fact that the mechanisms by means of which the solar wind torque is transmitted throughout this region are at present not well understood. Evidence derived from the analysis of helioseismological data suggests that within that portion of the radiative interior bounded by $r \approx 0.4\mathrm{R}_\odot$ and the base of the convection zone ($r \approx 0.7\mathrm{R}_\odot$), the solar angular velocity is only weakly dependent on depth and latitude, with a value about equal to that of the photosphere (see, e.g., Gough 1991, and references therein). Recent measurements of the rotational splittings of p-modes having low angular degree (i.e., $\ell = 1, 2$) further suggest that at great depth the angular velocity increases, attaining a mean value over the region $0.0 \lesssim (r/\mathrm{R}_\odot) \lesssim 0.2$ some 4.6 times that of the surface (Toutain and Frölich 1992). Note that if near-uniform rotation proves to be an accurate characterization of the rotational state of the bulk of the solar radiative interior, then angular momentum redistribution therein must be efficient, occurring over a time interval shorter than the evolutionary or spin-down time scales.

Such a state of affairs may not have prevailed early in the Sun's MS lifetime. Observations of rotation among solar-type stars in young, open clusters indicate that shortly after their arrival on the MS, many such stars sustain appreciable angular momentum loss (see, e.g., Stauffer 1991, and references therein). Indeed, from analyses of rotational velocity data for the most rapidly rotating G dwarf stars in the α Persei and Pleiades clusters, deceleration from velocities in excess of 50 km s^{-1} to values < 20 km s^{-1} in a time $\sim$ few $\times 10^7$ years is implied. As in the case of the Sun, this spin-down is presumably a consequence of the torque exerted on an individual star by the magnetically coupled wind it emits. However, the comparatively short rotational braking time inferred from observations is indicative of an enhanced rate of angular momentum loss; $\dot{J} \sim 10^4\ \dot{J}_\odot$ for rigid rotation, while $\dot{J} \sim 10^3\ \dot{J}_\odot$ if only the outer, convective envelope is braked. The latter, smaller value for $\dot{J}$ is more easily accounted for within the framework of standard, hydromagnetic stellar wind theories (Charbonneau 1992). This suggests that for stars of solar mass, the processes responsible for angular momentum transport from the core to the convection zone may operate with a time scale which is longer than the braking time, during the earliest phases of MS evolution.

Most extant computational models for the evolution of the Sun's internal and surface rotation assume rigid rotation of the convection zone, and treat angular momentum redistribution in the radiative interior by hydrodynamic means only (Endal and Sofia 1981; Pinsonneault *et al.* 1989; Tassoul and Tassoul 1989). Among the processes which have been utilized in studies of angular momentum transport in the non-convective portion of the solar interior are meridional circulation, Ekman flow, waves, and turbulent diffusion associated with the development of shear flow instabilities (Tassoul 1978; Endal and Sofia 1978; Zahn 1983). An alternative possibility is that angular momentum is redistributed inside the Sun by hydromagnetic means (see, e.g., Mestel and Weiss 1987; Roxburgh 1991). In particular, given the high electrical conductivity of the fluid of which the solar radiative interior is comprised, it seems reasonable to consider the consequences of the presence of a large-scale, primordial magnetic field $\mathbf{B}$ within that volume. From the induction equation, it is apparent that the local interaction between the velocity shear arising from wind-induced, differential rotation and the poloidal component $\mathbf{B}_p(= [B_r, B_\theta]$ in spherical polar coordinates) of such a field leads to the production of a toroidal magnetic field component B_ϕ. Associated with the growth of B_ϕ is the development of an azimuthally-directed Lorentz force which acts in such a way as to restore the sheared fluid to a state of uniform rotation. Note that the impulsive application of a transverse shear to a particular fluid volume element threaded by $\mathbf{B}_p$ gives rise to a field B_ϕ which propagates along $\mathbf{B}_p$ at the local Alfvén speed $u_A(= B_p/\sqrt{4\pi\rho})$. As a result of the Lorentz force exerted on each fluid parcel traversed by the disturbance, angular momentum is effectively redistributed along magnetic lines of force. Numerical simulations of the impulsive spin-up of a viscous, electrically conducting fluid indicate that for most poloidal field configurations, a state of nearly uniform rotation prevails within a time which is short compared to evolutionary time scales for the Sun or solar-type stars (Charbonneau and MacGregor 1992a; see also Moss, Mestel, and Tayler 1990).

Elsewhere in this volume, Charbonneau and MacGregor (1992c) present selected results from calculations performed as part of an effort to trace the rotational evolution of an internally magnetized 1 $M_\odot$ star, from the time of its arrival on the zero-age MS to the age of the present-day Sun. These computations pertain to a spherical star containing an axisymmetric poloidal magnetic field whose axis is aligned with the stellar axis of rotation. The convective envelope is taken to rotate rigidly at all times, and within the radiative interior the fluid velocity $\mathbf{u}$ and magnetic field $\mathbf{B}$ are assumed to have the forms $\mathbf{u} = u_\phi(r,\theta,t)\hat{e}_\phi$ and $\mathbf{B} = \mathbf{B}_p(r,\theta) + B_\phi(r,\theta,t)\hat{e}_\phi$. Starting from an initial state of uniform rotation, angular momentum is continuously removed from the convection zone at a rate in accord with the MHD wind model of Weber and Davis (1967). The evolution of u_ϕ and B_ϕ in response to the shear imposed at the periphery of the radiative interior is then calculated by numerical solution of the azimuthal components of the time-dependent induction and momentum equations. The reader is referred to the aforementioned paper by Charbonneau and MacGregor for additional details concerning the computational method and results.

In Figure 2, we show a sample of results relevant to the coupling between the convective envelope and radiative interior for a model with initial angular velocity Ω fifty times that of the present-day Sun, and in which the poloidal field $\mathbf{B}_p$ has dipolar symmetry with average strength in the range 0.1 to 10 G. For the column of panels labelled D2, $\mathbf{B}_p$ has a non-vanishing component B_r normal to the interface $r = R_{conv}$ over a restricted range of θ, while for the panels labelled D3, $B_r = 0$ at $r = R_{conv}$ for all θ. Depicted in the figure are: the induced toroidal field B_ϕ, at $r = R_{conv}$, $\theta = 45°$ (panels [A] and [B]); the torque (dJ/dt) applied to the convection zone by the wind, and the total magnetic (τ_B) and viscous (τ_ν) stresses exerted on the surface at $r = R_{conv}$ (panels [C] and [D]); a global measure of the differential rotation $\Delta\Omega$ (cf. Charbonneau and MacGregor 1992c) within the magnetized portion of the radiative interior (panels [E] and [F]). All quantities are shown as functions of time after the time of arrival on the zero-age MS. Note that for both the D2 and D3 configurations, B_ϕ exhibits an initial phase of approximately linear growth as time increases. During this period, the torque applied to the outer stellar layers by the wind is, for the most part, unbalanced by either the magnetic (in the D2 case) or viscous (in the D3 case) stresses which couple the envelope to the core. As is apparent from panels (E) and (F), the resulting spin-down of the convection zone is accompanied by an increase in $\Delta\Omega$, a quantity where instantaneous magnitude is proportional to the difference between the volume-averaged angular velocity of the radiative interior and the angular velocity of the envelope (cf. Charbonneau and MacGregor 1992c).

At later times, the rotational evolution of the envelope for both the D2 and D3 configurations is controlled by a near balance between the rate of angular momentum removal by the wind, and the rate of angular momentum resupply from the underlying radiative interior through the action of magnetic and/or viscous stresses. Of particular interest in this regard is the case in which the core and convection zone are directly coupled magnetically (e.g., D2). Inspection of panels (A) and (C) reveals that at the time when such an approximate balance of forces commences, the magnitude of B_ϕ is larger for smaller assumed B_p. This behavior can be understood by noting that when viscous and resistive dissipation are both negligible, the azimuthal component of the momentum equation is simply,

$$\rho \frac{\partial}{\partial t} \left(\Omega r^2 \sin^2\theta\right) = \frac{1}{4\pi} \, \mathbf{B}_p \cdot \nabla \left(r \sin\theta \, B_\phi\right) , \tag{10}$$

where we have written $u_\phi(r,\theta,t) = \Omega(r,\theta,t) \, r \sin\theta$. By integrating equation (10) over the core volume and applying Gauss' theorem, it follows that,

$$\frac{\partial}{\partial t} J_{core} = \frac{\pi}{4} R_{conv}^3 \; < B_r B_\phi > , \tag{11}$$

where J_{core} is the total angular momentum of the radiative interior and $< B_r B_\phi >$ is a weighted average with respect to θ at $r = R_{conv}$. The right-hand side of equation (11) represents the rate at which angular momentum is transported from the core to the envelope by the magnetic stress exerted at the interface. It is this quantity which, at later times, balances the torque applied to the envelope by the wind. Because for the solutions depicted the magnitude of dJ/dt does not depend strongly on the strength or geometry of the internal poloidal field (cf. panels [C] and [D]), the dynamical balance

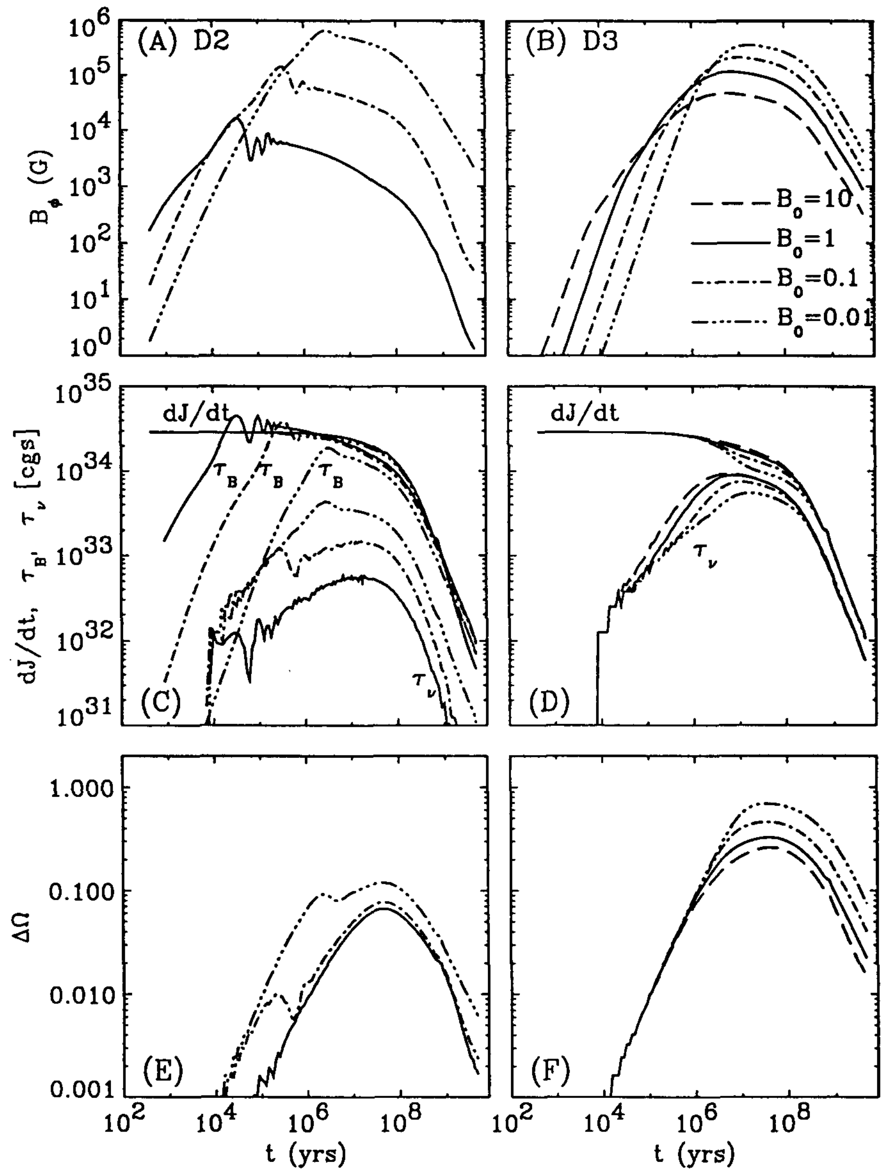

Figure 2: MS rotational evolution of an internally magnetized, 1 $M_\odot$ star as described in the text.

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$R^3_{conv} < B_r B_\phi > \approx \dot{J}_{conv}$ requires a larger value of B_ϕ for a smaller assumed value of B_p.

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