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Solar Cycle Variation of Anomalous and Galactic Cosmic Rays: the Role of Drifts and Current Sheet

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We apply a 2D Hoop-model to simulate the solar-cycle variation of anomalous and galactic cosmic rays (ACRs and GCRs) in a full 22-year cycle. Our axially symmetric model is 2D, but still captures the waviness of the heliospheric currents sheet (HCS). The strength of the heliospheric magnetic field (HMF), the corresponding diffusion coefficients and the waviness of the HCS are changed continuously. Special attention is given to the effects of particle drifts, some theoretical and computational aspects are discussed, including the handling and modifying the magnetic field in the polar regions of the Heliosphere. We briefly address the apparent disparity of ACRs and GCRs in the recent weak solar cycles when GCRs reached record high intensities while ACRs remained well below their previous intensity maxima. We find that, in addition to the faster diffusion, the weaker HMF and lower equator-to-pole potential, in itself, can also significantly lower the maximum energy ACRs can gain at the termination shock (TS). We also present numerical simulation for the 22-year variation of ACR fluxes at the TS at different heliographic latitudes.

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1. Introduction

The intensity of galactic and anomalous cosmic rays (GCRs and ACRs) track each other quite closely. The recent weak solar minima did, however, show an interesting disparity: while GCRs reached a record high intensity in 2009 [1], [2], the flux of ACRs remained below their level 22 years prior to this [3], i.e. during the previous solar minimum of the same polarity state of the Sun and heliospheric magnetic field (HMF). Since the transport of GCRs and ACRs are similar, the disparity suggested that the source of ACRs at the termination shock (TS) was lower, ie the acceleration of ACRs was less effective during this unusually weak solar minimum at 2009. Moraal and Stoker [4] were the first to point out that the faster diffusion should also result in lower rate of acceleration and a lower cutoff in the energy spectrum of ACRs at the TS. This interpretation has also been confirmed by recent numerical simulations [5]. While this mechanism is plausible and is most likely at work, the less efficient acceleration may also be influenced by a combination of other conditions: namely the weaker magnetic field [6] could in itself cause less efficient acceleration even if diffusion coefficients were to remain unchanged. In this work we consider the effect of the weaker magnetic field that corresponds to a lower pole-to-equator electrostatic potential.

We also investigate how the wavy current sheet (HCS) may interact with the TS. [7] found a marked difference in the Voyager ACR data in sectored and unsectored unipolar regions in the IHS beyond the TS. So this work is intended as a step toward resolving this question. In our simulation we shall adopt the paradigm, that ACRs are accelerated by diffusive shock acceleration at the TS [8] and apply our 2D 'hoop model' which, we believe, captures the most important features of the 3D HCSs. Our simulations are still preliminary, results are not intended to give quantitative fits to observations.

2. Model simulations and Discussions

Instead of the phase space density $f((x_i, p, t)$ we use $F = f p^3$ whose time evolution in momentum, p and space. x_i is described by an equivalent form of the diffusive Parker equation [9]

$$\frac{\partial F}{\partial t} = \frac{\partial}{\partial x_i}\left(\kappa_{ij}\frac{\partial F}{\partial x_j} - V_i F\right) + \frac{\partial V_i}{\partial x_i}\frac{\partial F}{\partial ln\, p^3} + Q \tag{1}$$

The consecutive terms describe anisotropic diffusion, advection by the solar wind at velocity V_i, and energy gain or loss due to the compression/expansion of the solar wind plasma. Diffusion is characterized by the 3D tensor, κ_{ij} which contains an anti-symmetric component, κ^A accounting for the regular spiral motion of cosmic rays is the global heliomagnetic field (HMF).

We adopt the paradigm, that ACRs are accelerated at the termination shock of the solar wind. Here we consider a spherical TS, we take typical values of 2.5 for the strength of the shock located at 90 AU from the Sun. We solve the Parker-equation on a fixed grid of heliospheric radius, r, polar angle, θ, and particle energy for singly ionized Oxygen from 0.010 to 100 Mev/nucl. A spherical free-escape boundary is assumed at r=120 AU.

In our calculation we assume azimuthal symmetry and employ a 2D 'Hoop' model, described in the next subsection.

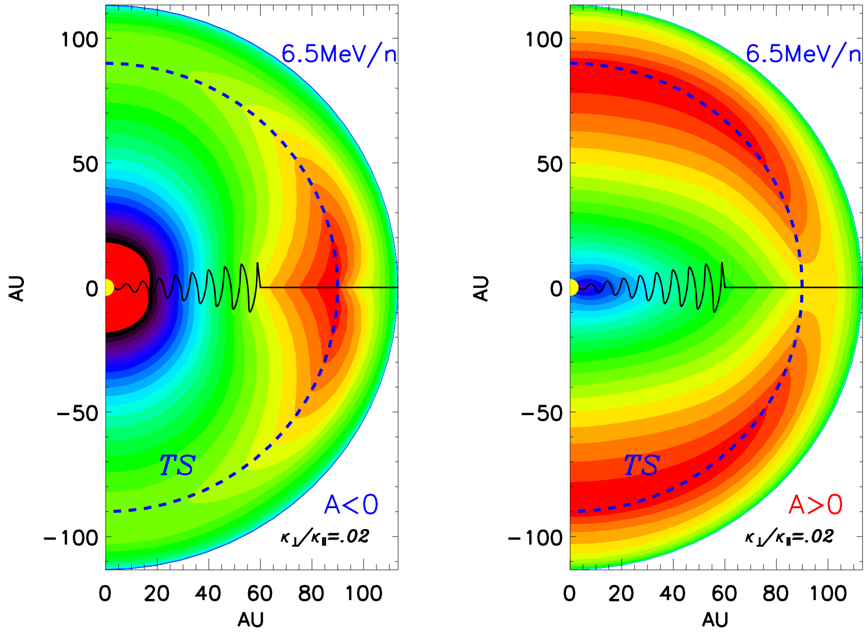

Figure 1: Sample contour-plots of ACR intensity in a kinematic hoop model during two consecutive solar minima. The contours show the well known polarity dependence in the meridional distribution of ACRs at solar minima. Black lines indicate the momentary HCS that is carried out radially by the solar wind.

2.1 Hoop Model of the Heliospheric Magnetic Field

This section describes our kinematic "hoop model" that has been applied for GCRs [10],[11]. The code includes acceleration at the TS and can readily be used for the acceleration and transport of ACRs. The virtue of this code is that this axially symmetric 2D code can still capture the most important 3D elements of the interaction of the wavy HCS with the TS.

The wavy HCS is an important element of cosmic rays transport. Particles can drift fast along the HCS, and the direction of drift reverses with the polarity change of the Sun. This waviness is an inherently 3-D feature whose precise incorporation requires 3-D codes, that are significantly more cumbersome than 2-D ones. 3-D codes are doable, but early 3-D codes [12],[13], had to make some compromise in resolution.

Our 'hoop model' is similar to that suggested by [14] to simplify the complex structure of the IHS. The Parker spiral of the HMF is tightly wound and in regions beyond 10 AU or so parallel diffusion acts mostly in the azimuthal direction and both the radial and latitudinal diffusion are dominated by the otherwise smaller perpendicular diffusion, $\kappa_\perp$. We construct a 'hoop' model, where the spiral field is substituted by expanding hoops that are carried out radially by the supersonic solar wind, that continues as an incompressible fluid beyond the TS. We note that the drift pattern remains essentially unchanged beyond 10 AU. This is also valid for the polar regions where the effect of fluctuating random transverse fields (see [15]) reduce the fast diffusion and drift that could otherwise prevail in the pure Parker spiral field.

The waviness of the HCS appears as time-dependence in the magnitude and polarity of the azimuthal B. We still use an anisotropic diffusion, $\kappa_{rr} = \kappa_\parallel cos^2(\psi) + \kappa_\perp sin^2(\psi)$ where $\kappa_\parallel$ and $\kappa_\perp$

stand for the parallel and perpendicular diffusion coefficients, while ψ is the garden-hose angle of the Parker spiral. The field is modified along the line of [15]. This modification is of particular importance at the polar regions. The problem of fast transport along the polar lines is somewhat obscure for GCRs, since the polar field lines are likely deflected and may never connect to the heliopause (HP). ACRs, on the other hand, are less sensitive to the outer boundary region (except for the escape of high energy ACRs) The next subsection will be devoted to some concepts of [15] and misconceptions [16].

2.2 Drift and Diffusion in the Polar Magnetic Field

Particle drifts are important, their importance increases with increasing rigidity. The original derivation of the 3D diffusion tensor by Parker [9] considered isotropic scattering, where the eigenvalues of the scattering operator are: $1/\tau_1, 1/\tau_2 + i\omega, 1/\tau_2 - i\omega$ with the parallel and perpendicular scattering times τ_1 and τ_2 assumed to be equal. At low rigidities, however, we can expect that particles forget their phase faster than their pitch angle, so τ_2 may be significantly smaller than τ_1, which could result in a larger $\kappa_\perp$ and a smaller κ^A. Drift effects are critically depending, both directly and indirectly, on $\kappa_\perp$ that could reduce drift effects [11].

The ideal Parker spiral of the heliospheric magnetic field (HMF) in the solar polar regions becomes extremely weak at larger radial distances, which would allow fast diffusion and drift through the polar areas. Jokipii & Kóta [15] pointed out that transverse field fluctuations, associated with the random walk of field lines and decreasing as $1/r$, may soon dominate the weakening radial polar magnetic field (weakening as $1/r^2$), and impede the fast transport of GCRs and ACRs along the polar lines. The inclusion of these fields increases the magnitude of the field, hence reduces the parallel diffusion and polar drift, too. Both diffusion and drifts depend inversely on the field strength. [15] adopted a substitution 'proxy' vector (D_i in what follows) as a plausible approximation to obtain more proper drift and diffusion coefficients.

This suggestion come under attacks recently in [16] and subsequent papers, arguing that the suggested 'enhanced magnitude' $D_i = (B^2 + db^2) \times B_i$ vector is not divergence free [16]. These authors miss the whole concept and do not realize the difference between a fixed vector and a fluctuating one. D_i is not and need not to be divergence free.

Indeed, $D_i = fB_i$ would be divergence free only if the multiplying factor f were constant along field lines, which would result in the same reduction at all radial distances. This not what [15] wanted and is not what fluctuating fields do. Their role do not remain constant but grow with increasing distance: initially small fluctuations become gradually dominating.

For a magnetic field, B_i, we have to take its inverse, B_i/B^2. For fluctuating field, the inverse of that may and likely will differ from the original field. In a somewhat mundane analogue the difference is resembling to the difference between real and complex numbers. Let us imagine that the field $B_i + i\delta B_i$ has an average field B_i which is known, and a fluctuating imaginary part $i\delta B_i$ whose average magnitude is known, but whose direction is random. Then the inverse is $(B_i + idB)$ is $(B_i - idB_i)/(B^2 + \delta B^2)$, that averages to $B_i/(B^2 + dB^2)$. In the process of averaging some features will be retained some others will be lost.

Another perhaps more relevant 'paradox': the average of a fluctuating unit vector is not a unit vector. This does not violate any law of physics or math.

Following standard procedures, the transverse fields can be inferred from a potential as $\delta B_i = \epsilon_{ijk} x_j \partial A/\partial x_k$ where $A(r, \theta, \phi)$ is a random function with a given mean square and spectrum. Then each single representation is automatically free of divergence. In a numerical code, however, we need the average drift and diffusion coefficients, which is obtained through nonlinear operations, typically involving the inverse of B_i.

While the statement of [15] may have been short, its meaning is quite plausible and self-explanatory. The 'proxy' vector, D_i was chosen so that it give a good approximation for the drift and diffusion coefficients, that is:

$$\frac{D_i}{D} = \left\langle \frac{B_i + \delta B_i)}{(B_i + B_i)^2} \right\rangle \approx \frac{B_i}{B^2 + \delta B^2} \tag{2}$$

The suggested $B_i/(B^2 + dB^)$ is a simple and plausible choice. The second part of Eq.(2) is not exact, a more precise result would depend on the structure of fluctuations [17]

Furthermore, [16] suggest that [15] would incorrectly continue at the transition through the TS. This is a mere guesswork of the authors, since [15] applied their simple formula to the old 1981 drift code of [18] which assumed a constant spherical wind and did not include the TS, so the question of transition was irrelevant and did not come up at all. The blind use of the given formula would indeed be an unfortunate oversight, and we were aware of this. It was an illustrative example and was not intended to be a universal rule. We may mention that most of the codes including the TS we developed that time adopted the condition that δB scales with the plasma density, n as $n(r)r$, which gives the correct transition.

Finally, [16] dismiss [15] as lacking physics base, which may be their personal opinion. On the other hand, however, we know that random fluctuations do exist, their presence was confirmed by Ulysses observations and were found to be consistent with the predictions of [19]. The inclusion of these fluctuating fields into modulation models has, in our view, quite strong physical base and motivation.

As a matter of subtlety: for the 'artificial' increase of the proxy substitute vector, D_i, a better choice is to increase its magnitude by the factor of $(1 + (\delta B/B)^2)$ rather than its square-root. The 'magnitude enhancement' can also be viewed as the inverse of the 'drift-reduction factor' obtained in several studies [17]. The agreement is not accidental. The message of [15], in simple words, is that, in calculating the diffusion and drift coefficients the expression in the denominator of B_i/B^2 should include the contribution of fluctuations.

The transverse fluctuations need not be limited to the polar regions, in our code they are applied at all latitudes and longitudes, their importance is, however, becoming gradually lower.

3. Model Simulations: ACRs at the Termination Shock

In this section we present numerical simulation results. First, we study the effect of the weaker HMF and corresponding lowe pole-to-equator potential. Second, we investigate the interaction of the wavy HCS with the acceleration at the TS. All our results are obtained for singly charged Oxygen densities at the TS at different latitudes.

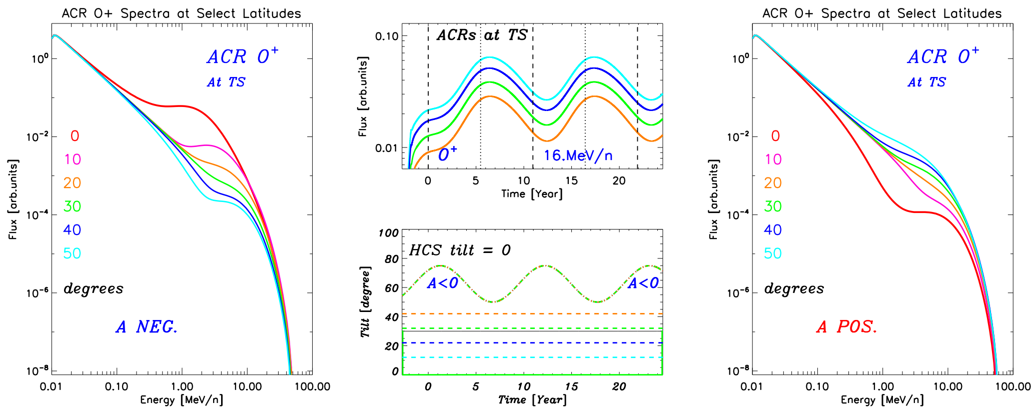

Figure 2: Variation of 16. MeV/n singly ionized Oxigens in response to a cyclic variation of the strength of the HMF with other parameters kept unchanged. The right panel shows the ACR spectrum at different latitudes. The simulation was obtained for A<0 negative cycle with flat current sheet. Similar results were obtained for A> cycle, are similar except that the latitudinal gradients are reverse

3.1 Role of Drifts and Pole-to-Equator Potential

The pole-to-equator electrostatic potential associated with the net magnetic flux of the rotating Sun plays an important role for both GCRs [20] and ACRs. Jokipii [21] pointed out that the maximum (total) energy of ACR ions should be in the range of the ion's charge times the electrostatic potential. In A>0 cycles ions drift from low to high latitudes in the acceleration process, so pick-up ions injected at or near the equator have the best chance to reach high energies. The opposite is true for A<0 negative cycles.

Our numerical calculation supports this expectation. Figure 2. depicts an A<0 scenario where the field strength is the sole parameter to change in and 11-year period by a factor of 1.5. ACR intensities follow the cyclic variation of B. Note that the figure show $1/B$. The right panel shows the energy spectrum at different latitudes. A distinct variation is present in medium energies, since drift effects increase with increasing energy but at the highest energies diffusion will diminish or fully erase the latitudinal gradients.

4. Role of Drifts and Heliospheric Current Sheet

The HCS is also an important element of the acceleration of ACRs. [22] showed that ACRs tend to be enhanced at the HCS during A<0 periods when they drift toward the HCS in both hemispheres during the acceleration process. The opposite holds for A>0, when ACRs are expected to be depleted at the HCS.

Figure 3. shows the simulated ACR fluxes in a 22-year cycles using fictitious parameters and changing the tilt-angle only. Finally shown in Fig.4 is a simulation when all of the HCS tilt, the HMF field strength and the diffusion coefficient change in 22 year cycle.

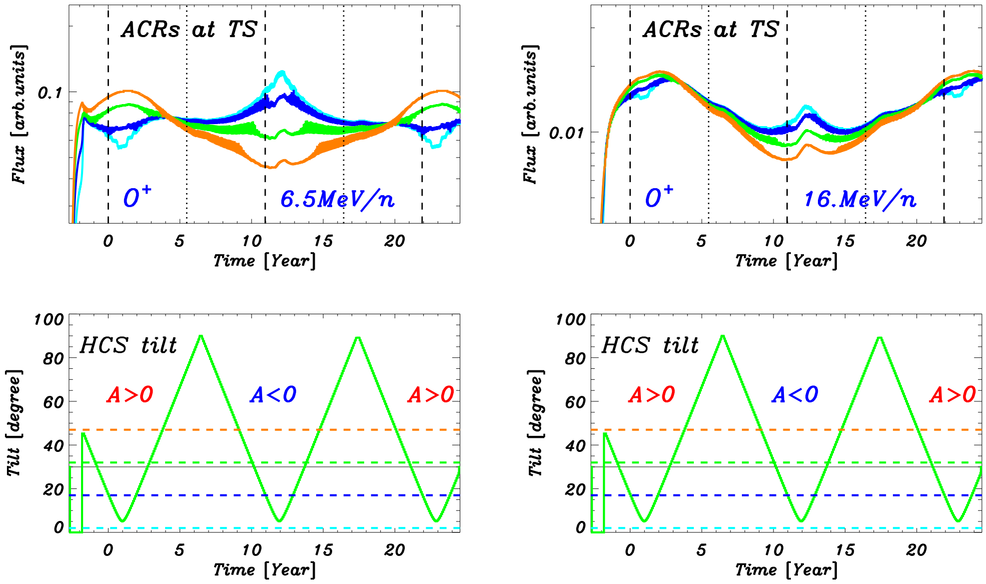

Figure 3: Variation of singly ionized Oxygen in response to a cyclic variation of the tilt of the HCS. Other parameters were kept unchanged. The panels show the ACR flux at different latitudes. Results obtained at higher energies show larger variation. Note the change of latitudinal gradients at the reversals of the polarity state of the HMF around years 6 and 17.

5. Summary

In this work we presented numerical simulations exploring the respective role of the strength of the HMF, and the waviness of the HCS. Preliminary simulation results tend to confirm our expectation, and we find that the weaker field, in itself, can significantly contribute to the decrease of the acceleration rate of ACRs in the recent solar minima.

References

[1] R. Mewaldt, A.J. Davis, K.A. Lave, et al., *Astrophys. J.*, **723**, (2010) 1

[2] I.G. Usoskin, A. Gil, G.A. Kovaltsov et al., *J. Geophys. Res.*, **122**, (2017) 3875

[3] R.A. Leske, A.C. Cummings, R. Mewaldt, & E.C. Stone, *Space Sci. Rev.*, **176** (2013) 253

[4] H. Moraal & P.H. Stoker, *Geophys. Res. Lett*, **115**, (2010) 12109

[5] R.D. Strauss, R.A. Leske, and J.S. Rankin, *Astrophys. J.*, **944**, (2023) 11

[6] E.J. Smith & A. Balogh, *Geophys. Res. Lett.*, **35** (2008) L22103

[7] M.E. Hill, R.B. Decker, L.E. Brown, et al., *Astrophys. J.*, **781** (2014) 94

[8] M. Pesses, J.R. Jokipii, & D. Eichler, *Astrophys. J.*, **246** (1981) 85

[9] E.N. Parker, *Planet. Space Sci.*, **13** (1965) 9

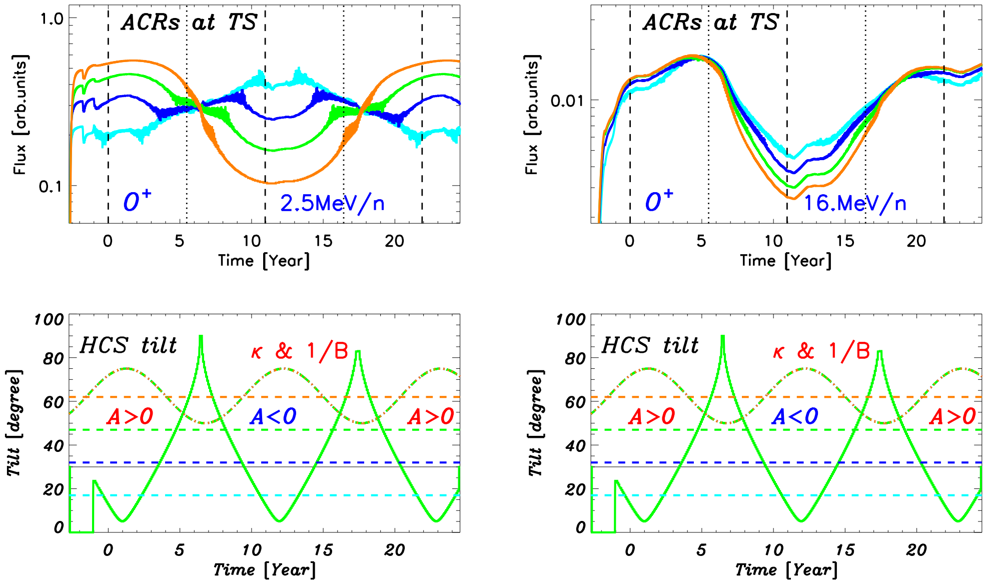

Figure 4: Same as Fig.3. except that the all three parameters change in a 22-year cycle. The sinusoidal curve in the bottom panel show the variation of κ and $1/B$ Note that the time-profile of the tilt angle is changed to increase the duration of flattish HCS.

[10] J. Kóta, *Space Sci. Rev.*, **176** (2013) 391

[11] J. Kóta, *JPhCS*, **767** (2016) id:012014

[12] J. Kota & J.R. Jokipii, *Astrophys. J.*. **265** (1983) 573

[13] J. Kota & J.R. Jokipii, *Geophys. Res. Lett.*, **18** (1991) 1797

[14] J.H. Drake, M. Swisdak, & M. Opher, *Astrophys. J.*. **709** (2015) 963

[15] J.R. Jokipii, & J. Kota, *Geophys. Res. Lett.*. **16** (1989) 1

[16] J.L. Raath, M.S. Potgieter, et al., *Adv. Space Res.*, bf 57, (2016) 1965

[17] N.E. Engelbrect, et al., *Astrophys. J.*, **841**, (2017) 107

[18] J.R. Jokipii, & J.M. Davila. *Astrophys. J.*, **248**, (1981) 1156

[19] J.R. Jokipii, J. Kóta, J. Giacalone, et al., *Geophys. Res. Lett.*, bf 23 , (1995) 3107

[20] J. Kóta, *Adv. Space Res.*, **1**, (1981) 135

[21] J.R. Jokipii, *Astrophys. J.*. **466**, (1996) 47

[22] J. Kota, & J.R. Jokipii, *Proc 23rd ICRC, Calgary*, (1993) 286