THE SNAKE: A RECONNECTING COIL IN A TWISTED MAGNETIC FLUX TUBE

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ABSTRACT

We propose that the curious Galactic center filament known as the Snake is a twisted giant magnetic flux tube, anchored in rotating molecular clouds. The MHD kink instability generates coils in the tube and subsequent magnetic reconnection injects relativistic electrons. Energy-dependent diffusion of electrons away from the coil produces a flat spectral index at large distances from it. Our fit to the existing data shows that the magnetic field of ~0.4 mG is large compared to the ambient ~7 µG field, indicating that the flux tube is force-free. If the relative level of turbulence in the Snake and the general interstellar medium are similar, then electrons have been diffusing in the Snake for about 3 × 10^7 yr, comparable to the timescale at which magnetic energy is annihilated in the major kink. Estimates of the magnetic field in the G359.19–0.05 molecular core are similar to our estimate of the magnetic field in the Snake, suggesting a strong connection between the physics of the anchoring molecular regions and the Snake. We suggest that the physical processes considered here may be relevant to other radio filaments near the Galactic center. We also suggest further observations of the Snake and other filaments that would be useful for obtaining further insights into the physics of these objects.

Subject headings: Galaxy: center — ISM: kinematics and dynamics — ISM: magnetic fields — MHD — radio continuum: ISM — stars: formation

On-line material: color figure

1. INTRODUCTION

The fundamental nature of the numerous radio filaments observed near the Galactic center (see LaRosa et al. 2000 for an overview) is unclear. In particular, how are they formed? What are the sources of the relativistic electrons? Why do they have such large magnetic field strengths? What is the reason for the unusually flat spectral indices in some of them? The discovery of the Snake (Gray et al. 1991), a 60 pc × 0.4 pc filament, approximately 150 pc to the west of Sagittarius A, resulted in the additional feature that the characteristics of the radio emission are strongly related to a morphological “kink.” Further observations (Gray et al. 1995) revealed a major and a minor kink, with the radio intensity and spectral index systematically varying away from the major kink. Gray et al. (1995) have summarized the various models that have been proposed to explain the spectacular elongated structure of the Snake. However, none of them have been able to successfully explain its major features. In our view, the Snake is in many respects one of the least complicated of the filamentary features close to the Galactic center, and a good physical model may hold the key to understanding the numerous arcs and filaments in the Galactic center region. We propose that the Snake is a magnetic flux tube with both ends anchored in dense rotating material (molecular clouds and/or associated H II regions). Differential rotation of the tube at either or both ends produces a monotonically increasing toroidal magnetic field, and when this reaches a critical value, “coils” or “loops” are formed through localized kink instabilities. Release of the magnetic energy stored in a coil, through magnetic reconnection (and possibly shocks), is the source of the relativistic electron energy. Energetic electrons diffuse away from each coil at an energy-dependent rate causing a flattening spectral index.

2. DETAILS OF THE MODEL

Let us take R to be the radius of the tube, \( L_{\text{tube}} \) its length, and \( B_\phi \) and \( B_z \) the toroidal and axial components of the magnetic field, in a cylindrical \((r, \phi, z)\) coordinate system coincident with the central axis of the unperturbed tube. As the ends of the tube are differentially rotated, \( B_\phi \) increases, and when \( B_\phi /B_z \sim 1 \), the kink instability produces a coil of magnetic flux with magnetic energy \( \Delta F_m \approx 4^{-1} \pi R^2 B^2 \) (Alfvén 1950, p. 117). We identify the observed major kink with such a coil; the minor kink may be another coil at a different stage of development.

Numerical simulations, in a solar physics context (Bazdenkov & Sato 1998; Amo et al. 1995), have shown that the coil’s magnetic energy is annihilated in a time on the order of a few transit Alfvén times, \( L_{\text{tube}} v_A \), where \( v_A \) is the Alfvén speed. Magnetic reconnection and possibly associated shocks in the coil can provide a source of energy for the acceleration of electrons to relativistic radio-emitting energies. Once accelerated, the electrons diffuse, and we model the diffusion of particles away from the acceleration region by a one-dimensional diffusion equation, assuming a uniform cross section for the tube. Taking \( f(p) \) to be the electron phase-space density, \( x \) the spatial distance along the flux tube, \( K(p) \) the diffusion coefficient, and \( C(p) \) the rate per unit volume of momentum space at which electrons are injected into the coil by the acceleration...
process, we adopt the diffusion equation

\[
\frac{\partial f(p, x, t)}{\partial t} - \frac{\partial}{\partial x} \left[ K(p) \frac{\partial f(p, x, t)}{\partial x} \right] = C(p) \delta(x).
\]

The delta function indicates that we are treating the coil, located at \( x = 0 \), as an infinitesimally small volume with respect to the rest of the tube. We also assume that the timescale for the dissipation of energy in the coil is greater than the time over which the electrons diffuse. We discuss this assumption further below.

We assume that \( K(p) = K_0(p/p_0)^n \), where \( K_0 \) is the diffusion coefficient at an electron momentum, of a GeV/c, and also take \( C(p) = C_0(p/p_0)^{1/2} \). Integrating equation (1) from \( x = 0^- \) to \( x = 0^+ \) gives the boundary condition at \( x = 0 \), namely,

\[
\frac{\partial f(p, x, t)}{\partial x} \bigg|_{x=0} = -\frac{1}{2} \frac{C(p)}{K(p)}.
\]

We normalize the diffusion equation and the boundary condition by introducing the distance \( L = 15 \) pc between the major and minor kinks as a fiducial length together with the normalized variables \( \tau = K(p) t L^2 = (K_0 t L^2)(p/p_0)^n \), \( \xi = x/L \), and \( g(\xi, \tau) = 2K_0 C_0^{-1} L^{-1} (p/p_0)^{1/2} f(p, x, t) \). In these variables, the diffusion equation and the boundary condition become

\[
\frac{\partial g}{\partial \tau} = \frac{\partial^2 g}{\partial \xi^2}, \quad \text{with} \quad \frac{\partial g}{\partial \xi} \bigg|_{\xi=0} = -1,
\]

for which the solution is

\[
g(\xi, \tau) = \left( \frac{4n}{\pi} \right)^{1/2} e^{-\xi^2/4\tau} - \xi \left[ 1 - \text{erf} \left( \frac{\xi}{\sqrt{4\tau}} \right) \right].
\]

The number density of particles per unit Lorentz factor \( \gamma \) is

\[
N(\gamma, \xi, \tau) = N_0(\gamma^{-1} g(\xi, \tau)),
\]

where \( N_0 = 2\pi(m_c c^3/(\gamma R_0 K_0))^{1/2} \). \( \gamma_0 = (m_c c)^{-1} p_0 \), and \( a = s + \beta - 2 \). We use \( N(\gamma) \) to evaluate the angle-averaged synchrotron emissivity \( \langle j_\nu \rangle \) as a function of frequency \( \nu \) from

\[
\langle j_\nu \rangle = \left( \frac{3e^2}{4c} \right) N_0 \nu_0 (2\nu/3\nu_0)^{(1-a)/2} \int_0^\infty y^{(a-3)/2} g(\xi, \tau) \hat{F}(y) dy.
\]

Here \( \nu_p = eB/(2\pi m_c) \) and the angle-averaged single-electron emissivity \( \hat{F}(y) = \int_0^\infty (1 - y^2/\nu^2)^{5/2} K_0(\nu) d\nu \). Because of the \( \gamma \) dependence in \( g(\xi, \tau) \), the spectral index near the kink is not simply \((a - 1)/2 \) but \((s + \beta - 2 - 3)/2 = (a - 1)/2 - \beta/4 \).

For a uniform flux tube of radius \( R \) and angular diameter \( \Phi \), imaged with a circular Gaussian beam with standard deviation \( \sigma \), the angle-averaged flux density per beam is \( F_\nu = 4(2\pi/3)^{1/2} (\sigma \Phi) sK_0(\nu) \). For an upper cutoff of 10 GeV, the deviations of the fit from the model are2 \((1000 \pm 200) \) GHz, for which \( \sigma = 1.1 \), \( R = 0.22 \) pc, and \( \Phi = 9.4 \) are appropriate. We restrict the fit to the region between Galactic latitudes of \(-3^\circ\) and \(-17^\circ\), that is, clearly related to the major kink. The radio images and the spectral index plot (Gray et al. 1995, Fig. 12) indicate that the minor kink influences the flux density south of \( b = -17^\circ \). The parameters of the fit (see Fig. 1) are \( N_0 = 3.5 \times 10^{-5} \) cm\(^{-3} \), \( B = 0.37 \) mG, \( a = 2.14 \), \( \beta = 0.57 \), and \( r_0 = 0.46 \). The deviations of the fit from the model are most likely related to systematic effects, such as variation in the field strength and width of the tube, rather than the statistical uncertainty in the data. The magnetic field can be estimated from the data, since it is related to the spectral index. However, experimentation with the fit shows that the precision of the estimate of \( B \) is not much better than a factor of a few although the result that the magnetic field is much greater than the ambient interstellar value, \( \approx 7 \) \( \mu \)G (Gray et al. 1995), is robust. Therefore, the flux tube is force-free.

If we assume that the injected electron spectrum extends between momenta \( p_1 \) and \( p_2 \), the total electron energy in the flux tube can be estimated as follows. Let \( A \) be the tube’s cross-sectional area; then the rate \( P_t \) of relativistic electron energy injection into the tube is

\[
P_t = 4\pi AC_0 \int_0^{p_2} \nu C(p) dp = 4\pi AC_0 \nu \int_0^{p_2} \frac{p_0^2}{p_0^2 - (p_0/p_0)^{a+1}}\nu \right)^{1/2} \delta_a^\nu \nu d\nu.
\]

The total relativistic electron energy in the flux tube is therefore

\[
E_{e, \text{tot}} = P_t t \approx 6.3 \times 10^{44} (p_2/p_0)^{a+1}\text{ergs},
\]

using the expression for \( N_0 \) (following eq. [5]) and the derived parameters. Time enters through the dimensionless model parameter \( r_0 = K_0 t L^2 \). For an upper cutoff of 10 GeV, \( E_{e, \text{tot}} \approx 3.9 \times 10^{35} \) ergs, and for 100 GeV, \( E_{e, \text{tot}} \approx 1.1 \times 10^{36} \) ergs. The available magnetic energy stored in the coil, \( E_B = 3.5 \times 10^{48} \) ergs, exceeds the total energy in electrons, as required for consistency of the model, but not necessarily by a large factor, depending on the maximum energy to which electrons are accelerated. This is a consequence of the high value of the derived magnetic field. These comparisons also suggest that the diffusion of electrons has been taking place for a time comparable to the total time available to annihilate the coil, consistent with our not

Fig. 1.—Model fit to the 1446 MHz (filled squares) and 4790 MHz (open circles) flux densities (Gray et al. 1995) for the Snake. The model fit only applies to the region associated with the major kink. [See the electronic edition of the Journal for a color version of this figure.]

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observing the process at a special epoch and consistent with the assumption of injection over the diffusion timescale.

The energy density of electrons $\epsilon_e$ can be derived simply from equation (5) for $N(\gamma)$. Near the coil ($\zeta = 0$), $\epsilon_e = N_im_e c^2 \times (-\alpha + \beta/2 + 2)^{-1} \left[(p_2/p_0)_{\alpha}^{\alpha} + (p_2/p_0)^{\beta/2} - (p_2/p_0)^{\alpha+\beta/2} \right]$, where $N_i = 2\pi^{-1/2}N_0 n e^{-\gamma_0 c^2}$. The energy density is insensitive to the upper cutoff, and for $p_2 = 10$ GeV/c, $\epsilon_e \approx 6 \times 10^{-11}$ ergs/cm$^3$, an order of magnitude larger than the energy density of the interstellar medium (ISM).

3. The Spectrum of Hydromagnetic Turbulence

It is usually assumed that relativistic electrons resonantly scatter off a preexisting level of hydromagnetic turbulence since resonant Alfvén waves are damped rapidly in the warm ISM by ion-neutral collisions, although this constraint is more important at higher than GeV energies (e.g., Melrose 1982). Let the energy density per unit wave number $k$ of resonant Alfvén waves (either preexisting or self-excited) be $W(k) = W_0(k/k_0)^{-\gamma}$, where $k_0 = eB_0c/p_0$ corresponds to waves resonating with 1 GeV electrons. Then the spatial diffusion coefficient for relativistic electrons is $K(p) = 4\eta(\eta + 2)/9\pi(\pi/2)^{1/2} [W(k_0)/W(k)]$, where $W_0 = B_0^2/\pi$ and $k_0 = eB_0c/p_0$ is the resonant wave number (Melrose 1982). Numerically, $K(p) = 1.4 \times 10^{19}(\gamma + 2)/(B_0^2)(W_0/k_0W(k_0))^{-1} \text{ cm}^2 \text{s}^{-1}$. Hence, $\beta = 2 - \eta$ and $\eta \approx 1.43$ for the parameters of our model. Our value of $\beta \approx 0.57$ is close to the Ormes & Protheroe (1983) value of 0.7 derived from a cosmic-ray propagation model. However, more recent models adopt a Kolmogorov value, $\beta = 1/4$, combined with the effects of “minimal reacceleration” (e.g., Ptuskin et al. 1999).

We have yet to take account of this possible effect. In order to constrain $\beta$ more effectively, one needs to take into account not only minimal reacceleration but the other physical parameters in the problem such as the width and strength of the flux tube and time variability in the injection process.

The time since the coil started to inject electrons along the flux tube is of interest for comparison with other timescales. This is $t = (L/K_0)\tau_0 \approx 3.1 \times 10^3 K_0 \tau_0 \text{ yr}$. For cosmic rays, Ptuskin et al. (1999) take $K_0 \sim 10^{26} \text{ cm}^2 \text{s}^{-1}$, implying that $K_0W_0 \sim 10^{20}$. If the same relative level of turbulence exists in the Snake, then $K_0 \sim 10^{26} \text{ cm}^2 \text{s}^{-1}$ since $K_0 \propto B^{-1}$ and $t \sim 3 \times 10^7 \text{ yr}$. According to the numerical simulations of comparatively short, wound flux tubes (Bazdenkov & Sato 1998), a coil disappears “explosively” in one to three Alfvén transit times. For the Snake, this is $\approx (1.5-4.5) \times 10^7 (n/10 \text{ cm}^{-3}) \text{ yr}$ based on $L_{\text{tube}} \approx 60 \text{ pc}$ and a number density $n \sim 10^{20} \text{ cm}^{-3}$ for the ambient ISM (Gray et al. 1995). Explosive bursts recur approximately every five Alfvén times, i.e., approximately every 7.5 $\times 10^7 (n/10 \text{ cm}^{-3})^{1/2} \text{ yr}$. Thus, for $K_0 \sim 10^{26} \text{ cm}^2 \text{s}^{-1}$, the time over which the electrons have been diffusing is on the order of the timescales of the energy-releasing process. Again, we are not required to invoke a special epoch of observation, and our assumption of continuous injection over the diffusion timescale is justified.

4. The Origin of the Magnetic Field

We suggest that the magnetic flux tube is initially anchored at both ends in molecular clouds before the latter undergo contraction and initiate star formation. Since molecular clouds condense from the warm ISM, the cloud magnetic field at the precontraction phase exceeds the general ISM value. Hence, such a flux tube emerging into the ISM from a molecular cloud would expand. As one or both of the clouds contract and rotate, a significant toroidal field is produced, and the resulting magnetic curvature force draws in the flux tube, thereby increasing the poloidal flux density in the entire flux tube to the molecular cloud value and eventually leading to a force-free state that becomes unstable for the reasons discussed above.

There is good evidence for the anchoring of the Snake in molecular clouds or associated $H$ ii regions. Observations by Uchida et al. (1996) reveal the northern end of the Snake intersecting an $H$ ii region in a CO “hole” in the molecular cloud complex G359.19–0.05. The supernova remnant projected against the southern end (Gray et al. 1995) provides circumstantial evidence for molecular material in that region. It is also possible that the Snake is part of a giant loop, the other end of which is anchored in another region of the Galactic center.

The generation of toroidal field is inextricably linked to the evolution of magnetic field and angular velocity in contracting molecular clouds. Here we discuss some of the physics involved. Nevertheless, a complete analysis involves substantial issues in the physics of contracting molecular clouds that are beyond the scope of this Letter.

The generation of toroidal fields by contracting clouds is related to the radiation of torsional Alfvén waves (see Melset 1999, p. 453). Taking $\rho_{e0}$ to be the mass density of the background ISM, the angular velocity, $\Omega$, and toroidal field are governed by

$$\rho_{e0} \frac{\partial\Omega}{\partial t} = \frac{B_0}{4\pi} \frac{\partial B_0}{\partial z}, \quad \frac{\partial B_0}{\partial t} = rB_0 \frac{\partial\Omega}{\partial z}. \quad (7)$$

The angular velocity in the background medium satisfies

$$\frac{\partial^2 \Omega}{\partial t^2} = v^2_0 \frac{\partial^2 \Omega}{\partial z^2}. \quad (8)$$

As a consequence of these equations, the toroidal field in a tube that is anchored only in the cloud is $B_{z0} \approx -r(4\pi\rho_0)^{1/2} \Omega_0 \approx -1.5 \times 10^{-6} (n/10 \text{ cm}^{-3})^{1/2} (\Omega_0 / \text{km s}^{-1} \text{ pc}^{-1}) \text{ G}$, where $\Omega_0$ is the angular velocity of the cloud (Melset 1999). In order to generate $B_z \sim 4 \times 10^{-4} \text{ G}$, one requires an extraordinarily large $\Omega \sim 270 \text{ km s}^{-1} \text{ pc}^{-1}$.

Consider now a tube additionally anchored, in another molecular cloud. We now have $\Omega = \Omega_0$, at $z = 0$ and $\Omega = 0$ at $z = L_{\text{tube}}$. On a timescale long compared to the Alfvén time, the solution of equation (8) is $\Omega = \Omega_0(1 - z/L_{\text{tube}})$, implying from the second of equations (7) that

$$\frac{\partial B_0}{\partial t} = -B_0 \frac{\partial \Omega_0}{\partial z}. \quad (9)$$

If $\Omega_0$ is constant for a time $t$, then after the flux tube has been twisted through an angle $\Delta\phi_0 = \Omega_0 t$, $B_z = -r\Delta\phi_0 B_0 / L_{\text{tube}}$. This expression can be derived simply from flux freezing in a twisted flux tube, neglecting the propagation time of Alfvén waves, and is used by Alfvén (1950) in the theory of the instability of twisted flux tubes referred to above. The derivation here makes a clear connection with the torsional Alfvén waves involved in the standard theory of magnetic braking. On the basis of this model, $B_z \sim B_0$ would be attained if $\Omega t \sim L_{\text{tube}}/R \approx 300 (R/0.2 \text{ pc})^{-1} \text{ rad}$. This would be the case, for example, for a cloud rotating at an angular velocity of $30 \text{ km s}^{-1} \text{ pc}^{-1}$ for a period of $10^7 \text{ yr}$.

The above estimates of the toroidal field assume that the flux tube has a constant radius, and this is clearly not satisfied in a contracting cloud. Moreover, the poloidal field changes under the opposing effects of compression and ambipolar diffusion, and the evolution of the field in the external ISM is also im-

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portant. However, our main point is that the anchoring of the magnetic field at both ends substantially affects the generation of toroidal field and the requirements on the angular velocity.

A related issue involves that of subcritical or supercritical contraction of the anchoring cloud(s). In the former case, the initial magnetic field supports the cloud and contraction occurs as a result of ambipolar diffusion accompanied by magnetic braking. The latter prevents the cloud from approaching centrifugal equilibrium (e.g., Basu & Mouschovias 1995 and references therein), and the angular velocity remains low, except for the latest, rapid stages of contraction when the core becomes critical. On the other hand, in the supercritical case, the magnetic field cannot halt contraction, and if the cloud has an initial angular velocity, angular momentum conservation causes it to spin up until it reaches centrifugal equilibrium. Further contraction then results from magnetic braking (Mestel & Paris 1984). This case may offer the best prospect for twisting of the flux tube. However, detailed investigation is deferred to future work. Note also that Basu & Mouschovias (1995) began their simulations with slowly rotating clouds.

An independent estimate of the poloidal magnetic field is of interest since this is one of the key parameters in our model and also determines whether the molecular cloud contraction associated with the Snake is sub- or supercritical. We use parameters derived by Uchida et al. (1996), who estimate that the clouds adjacent to the CO hole have masses \( \sim 5 \times 10^4 M_\odot \) with radii \( R \sim 2 \) pc. If we assume that these clouds have formed through compression of spheroidal regions (with semi-axes \( R_0 \) and \( K R_n \)) of the warm ISM of density \( n_0 \sim 10 \) cm\(^{-3} \) and with an ambient magnetic field \( B_{\phi} \sim 7 \times 10^{-6} G \) (Gray et al. 1995), then flux conservation implies that

\[
B_\phi \approx \frac{B_{\phi 0} R_0^2}{R^2} = \frac{B_{\phi 0}}{R^2} \left( \frac{3}{4 \pi K} \frac{M}{\mu \nu n_0 m_p} \right)^{2/3} \approx 7 \times 10^{-4} \left( \frac{B_{\phi 0}}{7 \times 10^{-6} G} \right) \\
\times \left( \frac{R}{2 \text{ pc}} \right)^{-2} K^{-2/3} \left( \frac{M}{5 \times 10^4 M_\odot} \right)^{2/3} \left( \frac{n_0}{10 \text{ cm}^{-3}} \right)^{-2/3} G.
\]

(10)

Using expressions derived in Mestel (1999, p. 429 et seq.), the critical magnetic field \( B_{\phi crit} \approx (6 \times 10^{-5} \text{ to } 1 \times 10^{-4})(M/5 \times 10^4 M_\odot)(/R^2 \text{ pc})^2 G \), with the lowest (highest) value for a spherical (flattened) cloud. These estimates raise the prospect that the clouds in the G359.19–0.05 complex are subcritical unless the original ISM region is highly prolate (\( K \gg 1 \)). However, our main point is that these estimates of both the actual and critical fields are similar to our estimate of \( 4 \times 10^{-4} G \) from our model for the Snake. It is therefore possible that the flux density in the associated contracting molecular cloud is not greatly enhanced over its initial value as a result of ambipolar diffusion. Many of the molecular clouds in the vicinity may have fields of this magnitude, but it is not until they are twisted and "lit up" by reconnection induced by an instability that they become observable. Interestingly, the Basu & Mouschovias (1995) simulations show that the magnetic field outside the supercritical core is indeed time independent.

5. DISCUSSION

High-resolution images (Gray et al. 1995, Figs. 10 and 11) show that the Snake appears to be split in two near the major kink, similar to flux tubes in the late stages of the Bazdenkov & Sato (1998) simulations, providing additional support for our proposal. As well as being directly applicable to the Snake, our model may open up a number of appealing possibilities for the dynamics of other magnetic filaments in the ISM near the Galactic center. It is feasible, in an environment with such a large density of molecular clouds, that rotation of magnetic flux tubes threading their cores would lead to kink instabilities and reconnection. The flattening spectral index of the Snake depends on the differential diffusion of energetic particles, and this is such a well-known phenomenon in the regular ISM that its role in this model is unremarkable. The index of the momentum dependence of the diffusion parameter that we have inferred for the Snake is in the range of values used in models of cosmic-ray propagation. More detailed models raise the prospect of a better determination of this index as well as the other parameters. In particular, a more precise value of the index of the momentum dependence of the particle creation rate, which our current estimates place near the value associated with strong shocks, is of interest. The type of data that is essential for more detailed modeling include high-resolution images showing more clearly the structure of the flux tube over its entire length and flux densities at different frequencies, enabling one to better constrain the energy dependence of the diffusion parameter and the magnetic field. It remains to be shown that flux tubes in the Galactic center can be wound up to such an extent that they produce a significant toroidal field. If an anchoring cloud acquires an angular velocity that persists long enough that it rotates through 300 rad (∼50 revolutions), that would be sufficient. Another possibility is that the required winding could be produced in a contracting cloud, although this idea requires further investigation. In our discussion of the magnetic field in the Snake, we have concentrated exclusively on the molecular complex at the northern end of the Snake. Currently, however, little is known about the southern end, and the masses and rotation rates of molecular clouds in that region are certainly of interest.

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