# KINEMATIC MAGNETIC RECONNECTION AT 3D NULL POINTS

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## ABSTRACT

Magnetic reconnection at a three-dimensional null point is the natural extension of the familiar two-dimensional X-point reconnection. Three dimensional null points are found in abundance in the solar corona, where they are thought to be an important feature with regards to the coronal heating problem. A model is set up here for reconnection at a 3D magnetic null point, by solving the kinematic, steady, resistive MHD equations in its vicinity. A steady magnetic field is assumed, as well as the existence of a localised diffusion region surrounding the null point. Particular attention is focused on the way that the magnetic flux changes its connections as a result of the reconnection process. The nature of the flows and the reconnection is found to be fundamentally different depending on the orientation of the electric current.

Key words: Magnetic reconnection, magnetohydrodynamics, magnetic null points.

## 1. INTRODUCTION

Magnetic reconnection is a fundamental process in many areas of plasma physics, whereby the magnetic field ($\mathbf{B}$) becomes restructured. Our ideas on how this restructuring occurs come mostly from the well-studied case of reconnection in two dimensions. In two dimensions, reconnection occurs at hyperbolic null points of the magnetic field. However, in three dimensions, reconnection can occur either at a null point or in the absence of a null point (Schindler et al., 1988; Lau and Finn, 1990).

The nature of reconnection at a three-dimensional null point is of particular interest since it is the three-dimensional generalisation of an X-point. Three-dimensional null points are also of crucial importance in the topology and interaction of complex fields on the Sun. They are found in abundance in the corona (see eg. Brown and Priest, 2001; Longcope et al., 2003), where their associated separatrices and separators are thought to be likely candidates for sites of coronal heating (Longcope, 1996; Antiochos, 2002; Priest et al., 2002). There is also evidence that null point reconnection may act as a trigger for at least some solar flares (Fletcher et al., 2001).

The local structure of the magnetic field around a three-dimensional null point is shown in Fig. 1. The *skeleton* of the null point is made up of a pair of field lines directed into (or out of) the null from opposite directions, known as the *spine*, and a family of field lines which are directed out of (or into) the null lying in a surface, known as the *fan* plane (Priest and Titov, 1996). If the current is zero then the null point is known as *potential*, the spine

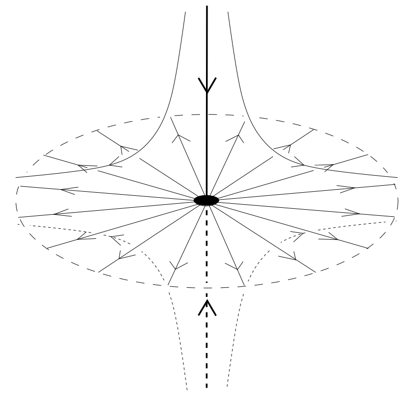

*Figure 1. The basic structure of a potential 3D null point.*

and fan are perpendicular, and fan field lines are purely radial. When a current component is present parallel to the spine, the field lines in the fan form a spiral structure, while a current component parallel with the fan plane results in the spine and fan 'folding' towards each other (see e.g. Parnell et al., 1996).

Priest et al., (2003) have shown that three-dimensional reconnection is fundamentally different in a number of important ways from the familiar 2D case. These details have been illustrated by a kinematic solution for reconnection in the absence of a magnetic null point (Hornig and Priest, 2003). Here we follow the same procedure for reconnection at 3D nulls. In Section 2 the differences between 2D and 3D reconnection are outlined. In Sections 3 and 4 the nature of reconnection is described at null points with spine- and fan-aligned currents. Finally, a summary is given in Section 5.

## 2. PROPERTIES OF 3D RECONNECTION

In order to investigate the evolution of magnetic flux in two dimensions, it is useful to define a flux transporting velocity $\mathbf{w}$ (Hornig and Schindler, 1996; Hornig and Priest, 2003) which satisfies

$$\mathbf{E} + \mathbf{w} \times \mathbf{B} = \mathbf{0}, \qquad (1)$$

which is possible in 2D since the electric field ($\mathbf{E}$) is always perpendicular to $\mathbf{B}$. Note that for reconnection to take place, the flux transporting velocity $\mathbf{w}$ must become singular at the null point, which is a signature of the breaking of the field lines (Hornig, 2001). By comparison with an ideal Ohm's law, we can consider $\mathbf{w}$ to be a flow within which the magnetic flux is frozen. The component of $\mathbf{w}$ perpendicular to $\mathbf{B}$ can be found from

$$\mathbf{w}_\perp = \frac{\mathbf{E} \times \mathbf{B}}{B^2}. \qquad (2)$$

However, in three dimensions this is not always possible. In general for 3D reconnection $\mathbf{E} \cdot \mathbf{B}$ is non-zero within a finite region $D$, with the result that in general no unique flux-conserving velocity exists, or in other words there is no unique field line velocity for field lines which thread $D$ (for a proof, see Hornig and Schindler (1996) or Priest et al. (2003)). We consider a finite region $D$ as the generic situation for the corona, since in general the magnetic Reynolds number is extemely high, and dissipation is likely to be enhanced only in well localised regions, for example when the presence of strong electric currents may drive micro-instabilities.

It is nonetheless possible to study the evolution of magnetic flux and field lines under certain circumstances. If no closed magnetic field lines exist within a localised non-ideal region ($D$) then we can still follow the motion of individual field lines from each end, since we know that in the ideal region on either side of $D$ they must remain attached to the same plasma elements for all time. Suppose the surface of $D$ is split into two parts, through one of which magnetic flux enters $D$, and through the other of which it leaves. Since each field line is anchored twice in the ideal environment, once on either side of the non-ideal region, one can follow the motion of the field lines by tracing them through space from either of these two sets of footpoints. By reference to the ideal environment then, it is possible to define a velocity with which the field lines passing *into* $D$ move, say $\mathbf{w}_{in}$, and another velocity $\mathbf{w}_{out}$ at which the field lines passing *out of* $D$ move. In the stationary case, these two velocities can be calculated throughout space by mapping the electric potential from each set of footpoints along the field lines. This leads us to the mathematical expressions

$$\mathbf{w}_{in/out} = \frac{-\nabla \Phi_{in/out} \times \mathbf{B}}{B^2}. \tag{3}$$

These two (pseudo-)field line velocities must each be identical to $\mathbf{v}_\perp$ on the relevant section of the surface of $D$, but in contrast to the two-dimensional case, they are not identical to each other inside $D$, nor are they identical to $\mathbf{v}_\perp$ on their continuations through $D$. This is a manifestation of the non-existence of a unique field line velocity, $\mathbf{w}$, as stated above. The result is that following field lines with $\mathbf{w}_{in}$ and $\mathbf{w}_{out}$, respectively, the field lines seem to split as soon as they are transported into the non-ideal region, and inside $D$ they continually change their connections.

The new properties of kinematic 3D reconnection described by Priest et al. (2003) are, briefly,

- The mapping of field line footpoints is in general continuous during reconnection, the exception being at null point separatrices.
- A unique field line velocity does not in general exist, so that in order to describe the evolution of the magnetic flux during reconnection it is necessary to use two field line velocities, as described above.
- Hence, field lines change their connections *continually* and *continuously* throughout the non-ideal region.
- Another way to look at this is to say that a field line traced through and beyond the non-ideal region moves beyond the non-ideal region in a virtual flow which is not the plasma velocity.
- There is no one-to-one reconnection of field lines, that is, for any field line on one side of $D$, there is in general no unique counterpart field line on the other side of $D$ with which its footpoints will become pairwise connected after the reconnection process.

## 3. SPINE-ALIGNED CURRENT

### 3.1 The Model

It is instructive to seek insight into the structure of the highly complex process of 3D reconnection by considering solutions to a reduced set of the MHD equations. We follow the method of Hornig and Priest (2003) and adopt the kinematic approximation, by solving the induction equation and Maxwell's equations, though not solving the equation of motion.

We seek a solution of the kinematic, steady, resistive MHD equations given by

$$\mathbf{E} + \mathbf{v} \times \mathbf{B} = \eta \mathbf{J}, \tag{4}$$

$$\nabla \times \mathbf{E} = \mathbf{0}, \tag{5}$$

$$\nabla \cdot \mathbf{B} = \mathbf{0}, \tag{6}$$

$$\nabla \times \mathbf{B} = \mu_0 \mathbf{J}. \tag{7}$$

The non-ideal term on the right-hand side of Eq. (4) is assumed to be localised. The nature of reconnection in the vicinity of a spiral null point, where $\mathbf{J}$ lies parallel to the spine, is investigated. In general the field in the vicinity of a spiral null point can be written (Parnell et al., 1996) in cylindrical polar coordinates as

$$\mathbf{B} = B_0 \left( R, \frac{jR}{2}, -2z \right). \tag{8}$$

where $B_0$ and $j$ are constants and $R = \sqrt{x^2 + y^2}$, such that $\mathbf{J} = (B_0 j/\mu_0)\,\hat{\mathbf{z}}$ is directed along the $z$-axis, which is coincident with the spine. Since the current is constant, it is necessary to localise the resistivity in order to obtain a localised non-ideal region. We take

$$\eta = \eta_0 \begin{cases} \left( \left(\frac{R}{a}\right)^2 - 1 \right)^2 \left( \left(\frac{z}{b}\right)^2 - 1 \right)^2 & R < a,\ z^2 < b^2 \\ 0 & \text{otherwise,} \end{cases} \tag{9}$$

where $\eta_0$, $a$ and $b$ are constant and $\eta_0$ is the value of the resistivity ($\eta$) at the null point. The null point and diffusion region are shown in Fig. 2.

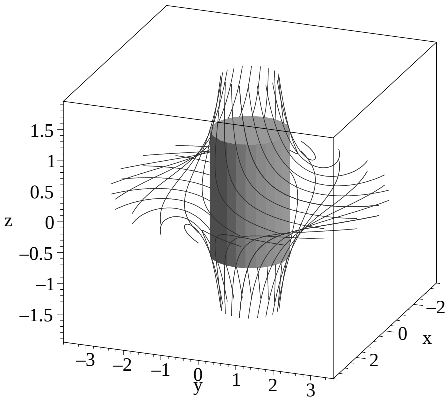


*Figure 2. Field lines on the boundary of the envelope enclosing the diffusion region (cylinder), showing the region of influence of the local solution (for $a = b = j = 1$).*

From Eq. (5), $\mathbf{E}$ can be written as the gradient of some scalar, $\Phi$ say, so that Equ. (4) becomes

$$-\nabla\Phi + \mathbf{v} \times \mathbf{B} = \eta\mathbf{j}. \qquad (10)$$

The component of this equation parallel to $\mathbf{B}$ may be integrated along field lines to give

$$\Phi = -\int \eta\, \mathbf{J}_{\|}\, dl + \Phi_0 \qquad (11)$$

$$= -\int \eta\, \mathbf{J} \cdot \mathbf{B}\, ds + \Phi_0(\mathbf{X_0}), \qquad (12)$$

where $s$ is a parameter along the field lines, and $\mathbf{X_0}$ are the field line footpoint positions. We choose to integrate $\Phi$ by setting $s = 0$ on surfaces $z = \pm z_0$, above and below $D$, and integrating in towards the $z = 0$ plane. Then the electric field ($\mathbf{E}$) and the component of the plasma velocity perpendicular to $\mathbf{B}$, $\mathbf{v}_\perp$, can be found from

$$\mathbf{E} = -\nabla\Phi \qquad (13)$$

and

$$\mathbf{v}_\perp = \frac{(\mathbf{E} - \eta\mathbf{J}) \times \mathbf{B}}{B^2}. \qquad (14)$$

When $\Phi_0$ is zero, an 'elementary solution' is found, upon which an ideal flow may be superimposed by choosing $\Phi_0$ non-zero. This is possible since, for a given magnetic field, Ohm's law (10) may be decomposed into a non-ideal component (15) and an ideal component (16) as follows

$$-\nabla\Phi_{non-id} + \mathbf{v}_{non-id} \times \mathbf{B} = \eta\mathbf{J} \qquad (15)$$

$$-\nabla\Phi_{id} + \mathbf{v}_{id} \times \mathbf{B} = \mathbf{0}. \qquad (16)$$

We have the freedom to add an ideal flow to the non-ideal solution since we do not solve the momentum balance equation here, which would otherwise determine the ideal part of the flow.

The mathematical expressions for $\Phi$, $\mathbf{E}$ and $\mathbf{v}_\perp$ are too lengthy to show here, but can be calculated in a straightforward way using a symbolic computation package. The full method of solution, as well as an in-depth analysis of the results, is described in Pontin et al. (2004a).

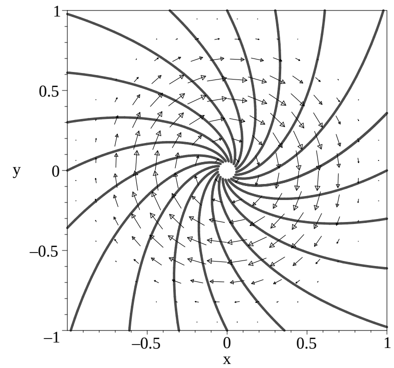


*Figure 3. Vectors of the plasma flow $\mathbf{v}$, along with a projection of the magnetic field lines, in the plane $z = 0.4$, for the parameters $B_0 = a = b = \eta_0 = 1$, $j = 2$.*

### 3.2 The Solution

We examine first the nature of the solution with $\Phi_{id} = \Phi_0 = 0$, i.e. just the local behaviour of the flux envelope enclosing $D$ with no extra ideal flow. Choosing to integrate Eq. (11) from $z = z_0$ we automatically start with $\Phi$ constant for $z > b$. Hence the electric field and plasma velocity are zero for $z > b$. The velocity for $z < b$ is a rotation within the flux envelope.

For the purposes of illustration of the results, it is at this stage convenient to add a component of $\mathbf{v}$ parallel to $\mathbf{B}$ such that $v_z = 0$. This is achieved by defining

$$\mathbf{v} = \mathbf{v}_\perp - \frac{v_{\perp z}}{B_z}\mathbf{B}. \qquad (17)$$

Our freedom to do this comes from the fact that Eq. (10) determines only the perpendicular component of $\mathbf{v}$, since it is the part that affects the behaviour of the magnetic flux, and so for our purposes the parallel component is arbitrary. The resulting flow pattern of $\mathbf{v}$ in a plane of constant $z$ is shown in Fig. 3. The radial component (as well as by definition the $z$-component) of $\mathbf{v}$ is zero, so we have purely rotational flow.

We can show that these rotational flows are not a result only of our particular choice of setup, for example the choice of $\eta$, but are in fact a consequence simply of the structure of $\mathbf{B}$ (and hence $\mathbf{J}$), and the presence of a localised non-ideal region. Consider the potential drop along sections of the closed loop illustrated in Fig. 4. $L1$ and $L3$ are radial lines in planes $z \geq b$ and $z < b$, respectively, while $L2$ is a field line on the surface of the envelope of flux threading $D$, and $L4$ lies along the spine of the null. The potential drop around any closed loop must be zero. The potential drop along lines $L1$ and $L2$ must be zero, as $L1$ lies at $z \geq b$ and $L2$ lies on the boundary of the envelope, and so $\Phi = \Phi_0$ all along them. Thus

$$\Delta\Phi_{L3} + \Delta\Phi_{L4} = 0,$$

or

$$\Delta\Phi_{L3} = -(\Phi_{spine} - \Phi_0) \neq 0, \qquad (18)$$

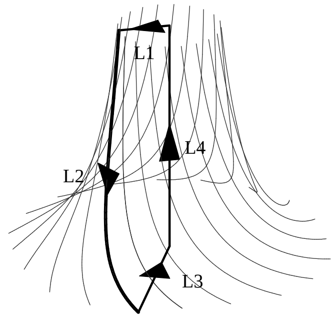


*Figure 4. Closed loop made up of line sections.*

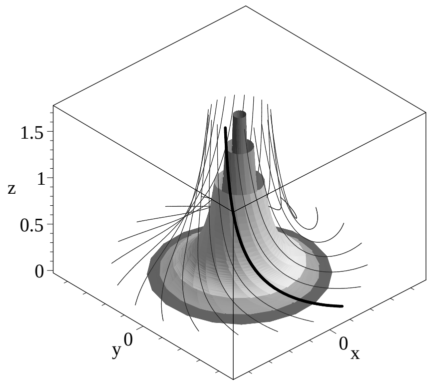


*Figure 5. Flux surfaces within which field lines periodically reconnect exactly with themselves.*

where $\Phi_{spine}$ is the value of $\Phi$ at the vertex of $L3$ and $L4$, which must be different from $\Phi_0$ since $\mathbf{E}$ is non-zero along the spine ($\mathbf{E}\cdot\mathbf{B} = \mathbf{J}\cdot\mathbf{B} \neq 0$ on spine). Since there is a potential drop along $L3$ there must also be a non-zero electric field along it. The electric field induces a plasma flow perpendicular to such a radial line, i.e. a rotational flow. It is found that this rotation has the same sense for $z > 0$ and $z < 0$, and has maximum magnitude in the $z = 0$ plane. Note that this argument is completely independent of the particular profile of $\eta$ within $D$.

Since the plasma velocity is rotational, the motion of field lines, and the nature of the reconnection is rotational as well. Thus field lines are continuously reconnected in 'shells' within the flux envelope enclosing $D$ (see Fig. 5). This means that in any arbitrary time, every field line within the flux envelope changes is connection (except the spine, due to symmetry). An initial field line which splits into two will be periodically exactly reconnected with itself, every time the half embedded in $R = a$ performs a full rotation of $2\pi$ radians. Note, however, that this period is different for each shell, since it is not a rigid rotation. For this reason, if we consider a flux tube within the envelope with a finite radial extent, there will be no periodic return to the initial state.

As discussed above, we may superimpose any ideal flow upon the 'elementary' solution described above. We would like to choose an ideal flow which shows the global effect of the local rotational slippage behaviour by transporting magnetic flux into and out of the local flux envelope. For this reason we choose to impose a stagnation-type ideal flow. Stagnation flows are also physically relevant flows to choose as they may perform the localisation of the physical quantities (Priest and Forbes, 2000; Hornig and Priest, 2003), and are common in reconnection solutions, such as Craig and Fabling (1996). We choose to set

$$\Phi_{id}(x_0, y_0) = \varphi_0\, x_0\, y_0, \tag{19}$$

on $z = \pm z_0$, where $\varphi_0$ is a constant. The resulting flow ($\mathbf{v}_{id}$) takes the form of a stagnation point flow, which is distorted due to the spiralling of the field lines.

The full plasma flow in this 'composite solution' is a competition between the rotational and stagnation flows. A number of different regions of flow exist, in each of which the rotational reconnection flow splits the field lines which are transported through the diffusion region. Two main types of behaviour are seen. In the first, flux tubes traced from either end seem to 'slip' apart relative to each other, but are carried away from the non-ideal region in the same direction (see pontin1.mpg). In another region of space however, a more classical type of reconnective behaviour is present, and initally joined flux tube footpoints are split by the reconnection, and end up connected to flux from completely different regions (see pontin2.mpg), with the crucial characteristic being that the separation of the two footpoints continues to increase in time. For a more detailed description of the process, see Pontin et al. (2004a).

## 4. FAN-ALIGNED CURRENT

The method of solution to model reconnection at a null point with fan-aligned current is exactly the same as described in the previous section. We choose

$$\mathbf{B} = B_0\,(x,\ y - jz,\ -2z)\,, \tag{20}$$

such that, without loss of generality, the current lies in the $x$-direction, and is given by $\mathbf{J} = (B_0 j/\mu_0)\,\hat{\mathbf{x}}$. The fan plane of this magnetic null point is coincident with the plane $z = 0$, while the spine is not perpendicular to this, but rather lies along $x = 0$, $y = jz/3$ (see Fig. 6). We this time prescribe a resistivity of the form

$$\eta = \eta_0 \begin{cases} \left(\left(\frac{R_1}{a}\right)^2 - 1\right)^2 \left(\left(\frac{z}{b}\right)^2 - 1\right)^2 & R_1 < a, z^2 < b^2 \\ 0 & \text{otherwise}, \end{cases} \tag{21}$$

where ${R_1}^2 = x^2 + \left(y - \frac{jz}{3}\right)^2$ and $\eta_0$, $a$ and $b$ are constants. $\eta_0$ is the value of $\eta$ at the null point, and the diffusion region is a tilted cylinder centred on the spine axis, extending to $z = \pm b$ and with radius $a$.

Performing the integration of Eq. (11) setting $s = 0$ on $z = z_0$ gives an expression for $\Phi(\mathbf{X}_0, s)$ for $z > 0$, and setting $s = 0$ on $z = -z_0$ gives $\Phi(\mathbf{X}_0, s)$ for $z < 0$. In order for these two expressions to match in the fan

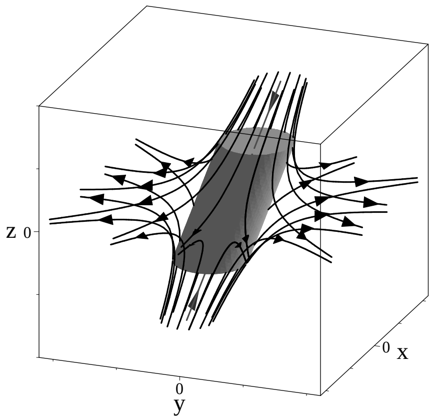


*Figure 6. 3D null point with fan-aligned current. The shaded cylinder shows the diffusion region.*

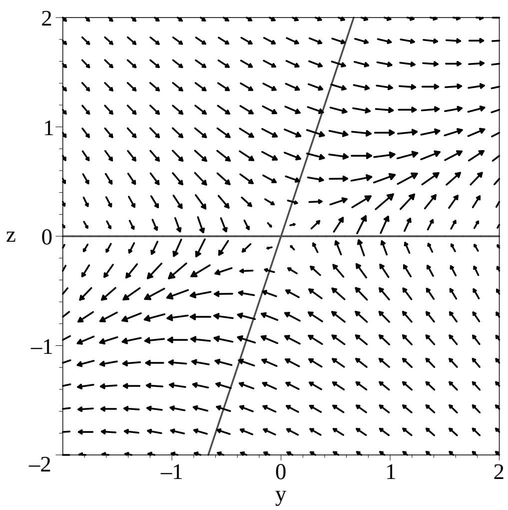


*Figure 7. The structure of the plasma flow, along with the fan and spine (black lines) in a typical plane of constant $x$, for parameters $\eta_0 = B_0 = j = a = b = 1$.*

plane, that is for $\Phi$ to be smooth and continuous, and thus physically acceptable, we must set $\Phi$ at $z = \pm z_0$ ($\Phi_0$) to

$$\Phi_0 = \frac{32}{21} \eta_0 B_0\, j\, x_0. \tag{22}$$

$\Phi\,(\mathbf{X})$, $\mathbf{E}$ and $\mathbf{v}_\perp$ can now be obtained as before.

In order to study the effect of the reconnection on the magnetic flux, we examine the plasma velocity ($\mathbf{v}_\perp$) perpendicular to $\mathbf{B}$. The nature of this flow in a plane of constant $x$ is shown in Fig. 7. The flow in the $x$-direction is zero across the spine, and is negligible for the reconnection process. The plasma flow is non-zero across both the spine and the fan, having a stagnation point structure in the $yz$-plane, centred on the null point. The result of the plasma flows across the spine and fan is that the nature of the field line behaviour under the reconnection is qualitatively the same as described by Priest and Titov (1996) in the ideal analysis, with field lines advected across the spine having a behaviour like their *fan reconnection*, and

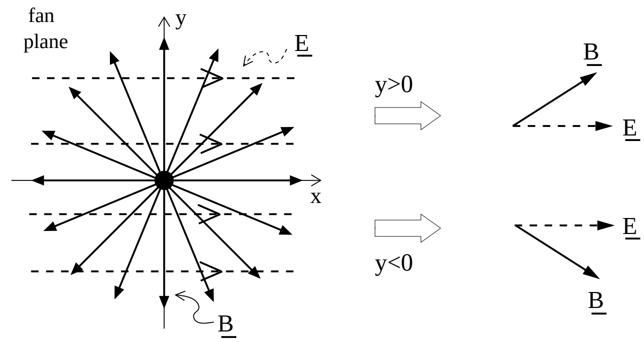


*Figure 8. The structure of the magnetic field ($\mathbf{B}$) and electric field ($\mathbf{E}$) in the fan plane. The non-zero vector product between the two implies a flow across the fan, in opposite directions for $y > 0$ and $y < 0$.*

those advected across the fan having a *spine reconnection* type behaviour.

Consider following field lines, with footpoints anchored in the ideal region, which pass across the top of the diffusion region ($D$) such that they pass down through $D$. These field lines can be seen to flip around the spine in the fan plane, with a behaviour analogous to the fan reconnection of Priest and Titov (1996) (see pontin3.mpg). By contrast, following field lines anchored in the flow across the fan plane we see these field lines move towards the fan until they lie in the fan plane, at which point they flip down the spine, and then move away from the null on the opposite side of the fan (as in spine reconnection).

Although a continuous stream of field lines is reconnected through the spine, no finite amount of flux is reconnected at it, since it is a single line. A finite amount of flux *is*, however, transported across the fan plane in an arbitrary time. This can be very important, since if we consider our null point to be surrounded by a global field, for example in the solar corona, then the fan plane is a separatrix surface of this magnetic field, which separates distinct regions of magnetic topology. The topology of this field changes when flux is transported across a separatrix surface. As described previously, in the solar atmosphere separatrices are thought to be very important locations for coronal heating.

The plasma flow across the fan plane is very important for topological considerations, and is a major difference between this solution and the one described in the previous section. It can be shown once again that the fundamental nature of the reconnection is simply a consequence of the structure of $\mathbf{B}$ and the presence of a localised non-ideal region. Consider the structure of $\mathbf{B}$ and $\mathbf{J}$, in the vicinity of the $x$-axis, in the fan plane. $\mathbf{J}\cdot\mathbf{B}$ has opposite sign for $x > 0$ and $x < 0$, since $\mathbf{J}$ and $\mathbf{B}$ are parallel on the $x$-axis for $x > 0$ and anti-parallel for $x < 0$. So $\Phi = -\int \eta \mathbf{J}\cdot\mathbf{B}\, ds$ has opposite sign for $x > 0$ and $x < 0$. Hence, in the fan plane close to the $x$-axis, $\mathbf{E} = E_x$ is uni-directional across the spine, as shown in Fig. 8. Since $\mathbf{E} \times \mathbf{B}$ is non-zero in the fan plane, we obtain a plasma flow in the $z$-direction, across the fan. Note that

$v_z$ must have different signs for $y > 0$ and $y < 0$ due to the different handedness of the vector products in these two regions (see Fig. 8). This argument is completely independent of the profile of $\eta$, and relies only upon the structure of $\mathbf{B}$ (and thus $\mathbf{J}$) and the fact that $\eta$ is localised. For a full description of the model and results, see Pontin et al. (2004b).

## 5. SUMMARY

A model for kinematic reconnection at 3D magnetic null points has been described. The behaviour of the magnetic flux in the process shows many properties which are very different from the familiar two-dimensional reconnection. The fact that no unique field line velocity exists results in a continual and continuous reconnection of field liens throughout the non-ideal region. It can be demonstrated that these new properties are still present in 3D reconnection when the full set of MHD equations are solved (Pontin et al., 2005).

It is found that the orientation of the electric current is crucial to the nature of the reconnection. When the current is parallel to the spine of the null point, a rotational type of reconnection is found within a flux envelope enclosing the diffusion region. However, when the current is parallel with the fan, magnetic flux is advected across both the spine and the fan of the null point. The fan-crossing flow is a particularly important result for the solar corona, where the fan surfaces of nulls separate regions of different magnetic topology. In order to change this topology, as the field must do to release some of the energy stored within it by photospheric motions, flux must pass across these fan surfaces.

## ACKNOWLEDGMENTS

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