Abstract.
Numerous variations have been proposed on the original suggestion by Blandford and Znajek that magnetic fields could be used to extract rotational energy from black holes. A new categorization of these variations is proposed so that they may be considered in a systematic way. “Black hole spindown” is defined more precisely, distinguishing decrease in the spin parameter \(a/M\) from decreases in angular momentum \(a\) and rotational kinetic energy, \(M - M_i\). Several key physical questions are raised: Can the “stretched horizon” of a black hole communicate with the outside world? Do accretion disks bring any net magnetic flux to the black holes at their centers? Is the magnetic field adjacent to a black hole force-free everywhere?

I BACKGROUND

Adding angular momentum to a black hole while keeping its mass fixed decreases its surface area:

\[
A = 8\pi M^2 \left[ 1 + \sqrt{1 - a_*^2} \right],
\]

where, as usual, \(M\) is the black hole mass, \(a_* \equiv a/M\) is its dimensionless spin parameter with \(a\) the specific angular momentum, and all quantities are in relativistic units (i.e., \(G = c = 1\), so that, for example, the unit of length is \(r_g = GM/c^2\)). Hawking and Ellis [12] showed that \(A\) cannot decrease; consequently, no matter what happens, a black hole of surface area \(A\) cannot ever have a mass less than \((A/16\pi)^{1/2}\), its “irreducible mass” \(M_i\). However, if \(a_* > 0\), \(M > M_i\). Therefore, reducing \(a_*\) can make \(M\) smaller; that is, braking a black hole’s spin can make it lose energy. In principle, this energy can be delivered in usable form to the outside world. If \(a_* = 1\), the energy potentially tappable is \((1 - 1/\sqrt{2})M \simeq 0.293M\). Comparing this quantity to the amount of energy released in the course of accretion (generally estimated as \(\sim 0.1M\)), we see that there is potentially as much energy stored in black hole spin as can be released in ordinary accretion.

The first proposal of a specific mechanism to extract this energy was due to Penrose [26]. His scheme made use of the fact that particles inside the ergosphere can achieve negative total energy (i.e., including their rest-mass energy) if their orbital frequencies are small enough (\(\Omega < a^{-1}(1 - r/2M)\)). However, Bardeen et al. [4] pointed out that it would be very difficult to accomplish this task by particle-particle events because there is an extremely large velocity difference between positive and negative energy orbits.

Since Penrose’s initial suggestion, a number of other ideas have been proposed. The unifying theme of all these proposals is to couple the black hole’s spin to

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external matter by magnetic forces. In some sense, they are non-local realizations of the Penrose process that make use of electromagnetic effects to convey negative energy into the hole, whether by injecting it with negative energy wave modes or negative energy particles.

Because the first such scheme was proposed by Blandford & Znajek [7], there has been a tendency in the literature to refer to them all generically as the “Blandford-Znajek” mechanism. Although all of them share the fundamental idea of magnetic couplings, and so in that sense are different versions of the same basic idea, there are now enough variations on the theme to make a clearer categorization useful. It is the object of this paper to present such a categorization, along with some comments about these variations’ relative standing.

II THE RATE OF ROTATIONAL ENERGY CHANGE

Before embarking on this discussion it is worthwhile to pause a moment to discuss several points of principle in order to clarify our language. First of all, let us list the ways by which black holes can gain or lose angular momentum (we will omit mechanisms involving singular events such as the initial creation of the black hole or a merger with a comparable mass object). One way is through the action of its own gravitational field: matter orbiting with an angular momentum vector oriented obliquely to the angular momentum of a black hole feels a torque. Here we will ignore this effect, assuming that any nearby matter orbits in the plane normal to the black hole spin. A second way is through matter crossing the event horizon bearing angular momentum. The third and final mechanism is through capture of photons with orbital angular momentum.

Next, let us consider what is required in order to identify the source of energy for a given event. It is occasionally assumed (e.g., in [19]) that all energy lost from stably-orbiting matter in an accretion disk is drawn from gravitational potential energy lost by accreting matter, with none coming from the spin of the black hole. However, this distinction—made locally—can be problematic because the shape of the gravitational potential depends on the black hole spin. To take an extreme example, when $a_\ast = a/M > 0.943$, the marginally stable orbit falls inside the ergosphere; in that case, part of the orbital kinetic energy of the disk material is due to the rotation of the black hole. If there are other ways of coupling the rotation of the black hole to the disk (see below), the identity of the source of energy at any particular location can become even fuzzier.

Although the local energy source can be ill-defined, it is always possible to make this distinction globally. All that is necessary is to compute the fluxes of angular momentum and energy across the black hole’s event horizon.

However, even for this distinction one must be careful about definitions. For example, if matter accretes with exactly the specific angular momentum $\hat{L}_{ms}M$ of the marginally stable orbit and the angular momentum of captured photons is neglected, the spin parameter $a_\ast$ always increases toward unity; when captured
radiation is included in the accounting, $a_*$ reaches equilibrium at $\simeq 0.998$ [33]. If the accreting matter arrives at the hole with a fraction $1 - \mathcal{L}_{ms}$ of $\dot{L}_{ms} M$, we may regard the spin as being tapped because the hole's rotational energy is less than it would have been if only the usual amount of energy (i.e. the binding energy at the marginally stable orbit) had been released.

One might also choose to impose the stronger condition that the hole is being “spun down” if $a_*$ actually decreases. Dividing the angular momentum fluxes into those depending on accretion of rest-mass and those depending on photon capture (photon capture also includes capture of electromagnetic waves), we find that the rate of change of $a_*$ is

$$\frac{da_*}{dt} = \frac{\dot{M}_o}{M} \left[ (1 - \mathcal{L}_{ms}) \dot{L}_{ms} - 2a_*(1 - \eta) \right] + \frac{\dot{M}_\gamma}{M} [J_\gamma - 2a_*],$$

where $\dot{M}_o$ is the rate of rest-mass accretion, $\eta$ is the fraction of rest-mass energy lost before matter arrives at the black hole, $\dot{M}_\gamma$ is the rate at which the black hole mass changes due to photon accretion, and $J_\gamma M$ is the mean angular momentum per photon. Note that $\dot{M}_\gamma$ depends on the energy of photons as measured at infinity; that can be quite different from the locally-measured photon energy. If photons are unimportant, whether $a_*$ increases or decreases depends on the balance between the angular momentum brought into the black hole $(1 - \mathcal{L}_{ms}) \dot{L}_{ms}$ and the angular momentum required to maintain the same ratio of spin to mass, $2a_*(1 - \eta)$.

It is apparent from equation 2 that $a_*$ can fall even while the angular momentum $a$ increases: the rate of change of $a$ is the same as for $a_*$ but without the two terms $\propto a_*$. Moreover, as equation 2 also shows, $a_*$ might remain constant even while $M$ increases. If that is so, the total rotational energy still increases. This fact suggests a still stronger definition—that the actual rotational energy $M - M_i$ decreases. Writing this criterion in terms of its component quantities, we have

$$\frac{d(M - M_i)}{dt} = \dot{M}_o \left\{ (1 - \eta) \left[ 1 - m_i - \frac{a_*^2}{2m_i(1-a_*^2)^{1/2}} \right] + \frac{a_* \dot{L}_{ms}(1 - \mathcal{L}_{ms})}{4m_i(1-a_*^2)^{1/2}} \right\}$$

$$+ \dot{M}_\gamma \left[ 1 - m_i - \frac{a_*^2}{2m_i(1-a_*^2)^{1/2}} + \frac{a_* J_\gamma}{4m_i(1-a_*^2)^{1/2}} \right],$$

where $m_i \equiv M_i/M$.

Thus, we see that the vague term “spindown” can connote any of several distinguishable results: diminishing $a_*$, $a$, or $M - M_i$. Because rotation contributes only a small amount to the total energy of the black hole when $a_*$ is small, the different criteria are very similar for $a_* \lesssim 0.8$; for larger values of $a_*$, the difference becomes more significant. When considering stored-energy reservoirs, the last definition is the most precise.
III  THE ORIGINAL BLANDFORD-ZNAJEK MECHANISM

As remarked above, the first plausible mechanism for removing black hole spin energy was proposed by Blandford & Znajek [7]. They imagined that as a black hole accretes, it would inevitably trap some net magnetic flux due to accreted field lines with connections to infinity. There would then be an approximately time-steady magnetic field configuration with field lines embedded in the hole’s event horizon, even while their far ends close at very large distance from the black hole. The enforced rotation of space-time due to the black hole spin would then drive an MHD wind. In their initial formulation, the field structure was supposed to be force-free everywhere, i.e. there was so little plasma attached to these field lines that $B^2/4\pi \gg \rho c^2$ everywhere. Phinney [27] extended this picture to include a small, but finite, plasma rest-mass density so that MHD wave group speeds would remain (slightly) less than $c$. Even in this modification, there is so little inertia outside the event horizon of the black hole that the orbital motion of nearby plasma has little effect on the rotation rate of magnetic field lines; only the distant region where the energy is delivered is considered to have any significant inertia.

Similarly, because the accretion rate is taken to be negligible, the change in the black hole’s rotational energy is due solely to the term proportional to $\dot{M}$ in equation 3; i.e., negative angular momentum and energy are brought to the black hole by zero rest-mass electromagnetic waves, not by ordinary matter.

These processes may also be visualized in terms of their associated electric fields. Suppose that the magnetic field is stationary with respect to a distant observer. Matter just outside the black hole is compelled to rotate with the hole, so in the matter’s frame there is an electric field perpendicular to both the field and the rotational velocity whose magnitude is proportional to the velocity mismatch between the field frame and the matter frame. This field drives a current from one field-line to another through the resistance of the horizon; the current then flows out along the new field-line until somewhere far from the black hole it crosses back to its original field-line and closes the circuit. When the resistance at infinity $R_l$ is comparable to the resistance in the horizon $R_h$, energy is dissipated in the distant “resistor” at a rate $\sim c\gamma g B_h^2$, where $B_h$ is the magnitude of the field at the event horizon. This impedance matching is equivalent to the condition that all the work done by the horizon forcing the field to rotate reaches infinity.

In the usual formulation, the magnetic field is taken to be time-independent and force-free everywhere between the black hole and a distant, but localized, “load”. A corollary to the assumption that the load is localized is that ideal MHD and flux-freezing apply everywhere between the black hole and the load. The force-free assumption then assures that wherever the field and the plasma move in lock-step the field controls the rotation rate. Given those assumptions, matching the current through the horizon to the current through the load determines the rotation rate of the flux-lines linked by the current loop:
\[ \Omega_F = \frac{\Omega_h R_l + \Omega_l R_h}{R_h + R_l}, \]  

(4)

where \( \Omega_h \) is the rotation rate of the black hole and \( \Omega_l \) the rotation rate of the load [5]. The load’s rotation rate is determined by a balance between whatever forces (gravitational, magnetic, . . .) act upon it.

When thinking in terms of the equivalent circuit, it is important to distinguish true dissipative resistance from effective resistance (this distinction is closely related to the distinction between ordinary and radiation resistance). For example, if the field-lines at infinity pass through plasma with high conductivity, the major part of the energy transmitted to infinity can be in the form of bulk work; i.e., the distant matter can be accelerated coherently. By contrast, ordinary resistance produces heat. The impedance matching referred to earlier refers to the total load at infinity, not solely the dissipative part.

Because \( B_h \) depends on the history of past accretion onto the black hole, specifically the total amount of net magnetic flux accreted, it is unclear how large \( B_h \) should be. Because there is resistance in the horizon, the field must be supported by currents somewhere in the vicinity of the black hole. Rees et al. [31] suggested that they might be located in a “fossil” accretion disk so that they might not be any farther from the horizon than a distance \( \sim r_g \). Rees et al. then argued that the field could not be any larger than \( \sim (r_{ms}/r_g)^2 \) times the pressure in the fossil disk, or there would be no dynamical equilibrium. Another way to set the scale is to suppose that there is a small amount of accretion, and estimate \( B_h \) as comparable to the field strength in the nearby accretion disk: if, as seems likely from recent numerical simulations [8,32], \( B^2 < 8\pi p \), this estimate yields a result somewhat smaller than the previous one [10]. Combining the more conservative estimate with the assumption of good impedance matching, the expected luminosity is

\[ L_{BZ} \sim \dot{M}_o c^2 (v_{orb,ms}/c)(r_{ms}/h_{ms})(r_g/r_{ms})^2, \]

(5)

where everything with subscript \( ms \) is evaluated at the marginally stable orbit.

There are several issues regarding this mechanism that remain incompletely understood. Punsly & Coroniti [29] raised the question of whether the black hole event horizon and plasma far from the black hole could be in causal contact. They argued that, although the plasma density is taken to be negligible in the usual formulation of the Blandford-Znajek mechanism, it cannot be literally zero. There would then be some accretion onto the black hole, and the fast magnetosonic speed would be (slightly) less than \( c \). If so, there would inevitably be a surface surrounding the event horizon within which the inward velocity of the plasma would be greater than the fast magnetosonic speed, and no signal carrying energy or information could propagate outward across that surface. Punsly [28] further argued that, because of severe gravitational redshifting, electromagnetic waves (in this case, the fast magnetosonic mode) can carry only tiny amounts of energy away from the event horizon. The significance of these criticisms is still unclear.
Further questions may also be raised about this form of the Blandford-Znajek scenario: Must the black hole necessarily accumulate a net flux? Perhaps all field lines suffer enough reconnection in the course of accretion that only closed loops are brought to the black hole. If this is the case, the field loops would close near the black hole, not far away. Would the mechanism still work if the plasma is not magnetically-dominated everywhere between the black hole and the load? Field lines passing through the accretion disk, for example, will surely have substantial inertia attached. Similarly, the load region (presumably near the Alfvénic surface of the wind where $B^2/4\pi \sim \rho v^2$) might not be too far from the black hole even if the field lines close at much greater distance. One immediate consequence of non-force-free behavior would be to change equation 4 to involve a mean value of $\Omega_l$ weighted inversely by the local resistance between the magnetic flux surfaces. Two other consequences would be more serious, however. Significant plasma inertia would cause the magnetosonic speed to fall well below $c$, so that the inner magnetosonic surface would likely move well outside the event horizon. Plasma inertia would also slow the approach to stationarity and possibly disrupt it altogether. In that case, the whole time-steady picture—a fixed magnetosonic surface, well-defined field rotation rates, a steady-state lumped-parameter electrical circuit analog—would be undercut.

Differing answers to these questions may be regarded as the basis for the variations on this scheme that have been suggested. Alternatively, they may be used as the structure of a classification scheme. We will organize the remainder of this paper along those lines, dividing schemes according to their answers to three questions:

- Does the innermost part of the magnetic field run through the event horizon or through plasma in the ergosphere?
- Do the field loops close nearby (e.g., in the disk) or far away?
- Is there anywhere in the system other than the horizon and a localized load where plasma inertia matters?

Labelled by this scheme, the original Blandford-Znajek mechanism is one in which the field lines run through the event horizon and out to infinity, and the field is force-free everywhere except in the event horizon and the load.

**IV FIELD LINES ANCHORED OUTSIDE THE HORIZON, CLOSED AT INFINITY, AND FORCE-FREE**

The first alternative was invented by Punsly and Coroniti [29,30] and elaborated by Punsly (1996). It has also been recently discussed by Li [18]. They suggested that a better means to extract the rotational energy of the black hole might be a magneto-centrifugal wind anchored in plasma orbiting in the ergosphere, but well away from the event horizon (see also [24] for a similar idea). This device would finesse the possible causality problem of field lines tied to the event horizon,
but would be otherwise quite similar to the original Blandford-Znajek idea: the
field lines extend to infinity, the structure is supposed at least roughly time-steady,
and the plasma inertia is required to be low enough as to satisfy the condition of
magnetic domination everywhere but possibly in the distant load.

Preliminary numerical simulations of a version of this idea exist [15]. In these
simulations, the initial magnetic field was taken to have uniform intensity and to be
directed parallel to the rotation axis. Although the Alfven speed is relatively small
in the equatorial plane, it is $\sim c$ everywhere outside the disk. As a result, magnetic
braking of orbiting material is very strong, leading to dramatic infall, shocks, and a
(transient?) jet that is pressure-driven along the axis but magnetically-propelled at
larger radius. Just as for the classic Blandford-Znajek mechanism, the characteristic
luminosity scale is $\sim |\vec{B}|^2 r_h^2 c$, but the coefficient could be considerably less than
unity. Based on this point of view, Meier et al. [21] and Meier [20] have also
suggested that the relative strength of the magnetic field at the base of the jet could
explain the morphological contrast between FR1 and FR2 radio galaxies.

V FIELD LINES ANCHORED IN THE HORIZON,
CLOSED IN THE DISK, AND FORCE-FREE

In the original Blandford-Znajek scheme, the field lines are anchored in the event
horizon, but extend to infinity. It is also possible that some or all the field lines
threading the horizon could instead close by passing through the disk. Because the
material in the disk will arrive at the black hole in an infall time, the details of
such a structure must be transitory. However, by the same token, if the loops close
well outside the inner edge of the disk, the characteristic duration of this structure
would be $\sim \alpha^{-1}(r/h)^2$ dynamical times, where $\alpha$ is the usual ratio of $r-\phi$ stress to
local pressure and $h$ is the disk thickness. Particularly in a thin disk, this could be
a relatively long time.

For this reason, some discussions of this scheme have assumed a time-steady
situation and its corollary, field lines with a fixed rotation frequency (e.g., [17]; but
see [11] for an intrinsically non-steady version). In this case, because the resistance
of disk plasma is so much less than $R_h$, $\Omega_F$ for all field-lines must be very nearly
the rotational frequency $\Omega_d$ of the disk matter they thread. To estimate the offset,
we observe that the $\vec{J} \times \vec{B}$ force can be written in two equivalent ways. Most often,
this force is evaluated by using Ampère’s Law (neglecting the displacement current)
to write $\vec{J} = (c/4\pi) \nabla \times \vec{B}$. However, one can also use Ohm’s Law to compute the
current, i.e. set $\vec{J} = \sigma(\vec{E} + \vec{v}/c \times \vec{B})$. Identifying the field-line velocity with the
$\vec{E} \times \vec{B}$ drift speed [14], we can then equate the two forms and find

$$|\Omega_F - \Omega_d| \sim \frac{c^2}{8\pi h r \sigma}$$

where $h$ is the disk thickness and $r$ is the radial coordinate of the place the field-
lines pass through. If resistivity in accretion disks is due to electron-ion Coulomb
collisions, $\sigma \sim 10^{14} \text{s}^{-1}$; even if particle-wave scattering or turbulence magnifies the effective collision rate, $|\Omega_F - \Omega_d| \ll \Omega_d$.

To compute the rate at which black hole rotational energy is tapped, one must next estimate the total resistive load due to the disk. The energy per unit volume dissipated by true electrical resistance is the characteristic rate of dissipation for turbulent magnetic field in the disk, $\sim (B^2/8\pi)\Omega$, times a very small factor $\sim c^2/(v_{\text{orb}}\sigma h)$. On the other hand, if $\langle B_r B_\phi \rangle \neq 0$, the rotational work done by the field is likely to be much larger. The problem is that the field structure within the disk is very much not force-free, so the usual simplifications invoked to estimate power output in the Blandford-Znajek mechanism don’t apply. As a result, detailed calculations are necessary to evaluate the true load—and therefore the power that is actually generated—in this model.

**VI FIELD LINES ANCHORED IN PLUNGING PLASMA, CLOSED IN THE DISK, AND NOT FORCE-FREE**

In contrast to the previous two pictures, let us now consider the consequences of a disk whose material is actually accreting. Within it, the growth of the magneto-rotational instability to nonlinear amplitude creates strong MHD turbulence. When plasma follows circular orbits whose frequency declines outward, the shear automatically biases the turbulence so that $\langle B_r B_\phi \rangle < 0$ and angular momentum is carried outward, thereby permitting accretion [3].

When material leaves the marginally stable orbit and plunges inward, its high electrical conductivity assures that it carries magnetic flux along. Whether or not there is any net magnetic flux threading the black hole’s event horizon, most magnetic field lines in the plunging plasma are closed relatively close by, for the MHD turbulence in the disk should lead to much reconnection. The question now arises as to whether these field lines, threading plasma plunging along unstable orbits through the ergosphere, can exert enough stress to tap the black hole’s rotational energy at an interesting rate.

Novikov & Thorne [23] argued that in a time-steady, axi-symmetric disk the stress $T_{r\phi}$ should approach zero at the marginally stable orbit because the inertia of material in the plunging region inside $r_{ms}$ would be too small to exert any significant stress. If this were the case, this version of the Blandford-Znajek mechanism (and the previous one!) would be unimportant. However, as remarked in [33,25], it is possible for magnetic forces to be significant even in the absence of substantial inertia. If so, $T_{r\phi}$ need not go to zero at $r_{ms}$.

In fact, the same MHD processes that account for angular momentum transport in the disk proper are likely to apply in the neighborhood of $r_{ms}$. The existence of the magneto-rotational instability depends only on the disk shear; it is therefore just as strong in the vicinity of $r_{ms}$ as it is farther out in the disk. There is also no particular reason why MHD turbulence should damp more quickly in the inner
disk than the outer. We may therefore expect the amplitude of MHD turbulence (normalized to the disk pressure) to remain roughly constant as \( r \to r_{\text{ms}} \). Because the shear is still strong, we can likewise expect \( T_{\phi}/T r(T) \) to not change dramatically in this region. In fact, if anything, the ratio of magnetic stress to pressure is likely to rise near \( r_{\text{ms}} \) because the pressure begins to decline as the inflow velocity rises, whereas we can expect the magnetic field strength to vary more slowly. It therefore appears that there is no local mechanism to enforce the decline in \( T_{\phi} \) demanded by the Novikov-Thorne model.

If this is the case, flux-freezing through the plunging region can lead to magnetic forces becoming comparable to gravity in a roughly time-steady and axi-symmetric accretion flow when \( v_r \sim v_\phi \) \[16\]. \( T_{\phi} \) at \( r_{\text{ms}} \) large enough to significantly alter the accretion efficiency follows as a corollary \[9\]; in a time-steady, axi-symmetric state, the additional dissipation per unit area in the disk falls roughly as \( r^{-7/2} \) at \( r \geq r_{\text{ms}} \) \[1\].

Recent simulations by Hawley & Krolik \[13\] (3-d non-relativistic MHD in the pseudo-Newtonian Paczyński-Wiita potential) support these arguments (but see also \[2\]). They find that the inner regions of disks become highly turbulent and non-steady, but that the azimuthally-averaged and vertically-integrated magnetic stress is essentially flat from the innermost part of the disk proper through the marginally stable orbit and well into the plunging region. One might therefore expect the accretion efficiency to be somewhat greater than the one predicted on the basis of zero-stress at \( r_{\text{ms}} \) even without black hole rotation.

Black hole rotation (and its attendant frame-dragging) may enhance this supplemental energy release because it is mediated by magnetic torques \[9\]. In other words, when the black hole rotates, this mechanism is very similar to the previous one, but transfers the inner field “anchor” to plunging plasma and does not require either force-free field structure or time-steadiness.

As for any mechanism proposed to tap black hole spin, “spindown” must be defined carefully. If the weakest of the three criteria is used, \( \text{any} \) increase in the accretion efficiency above the nominal is associated with a diminution of the black hole’s rotational energy, for it means that the accreting matter enters the hole with smaller angular momentum than it might otherwise have. If the constant \( a_\ast \) criterion is used and captured radiation is ignored, spindown occurs when the efficiency is greater than

\[
\eta_{\text{eq}} = 1 - \frac{\sqrt{1 - 3/r_{\text{ms}}} + 2a_\ast/r_{\text{ms}}^{3/2}}{1 - a_\ast/r_{\text{ms}}^{3/2}}; \tag{7}
\]

\[1\]. For \( 0 \leq a_\ast \lesssim 0.95 \), \( \eta_{\text{eq}} \simeq 0.3 - 0.35 \). Captured radiation diminishes \( \eta_{\text{eq}} \) by \( \simeq 0.1 \) for \( a_\ast \gtrsim 0.99 \). The contrast between this criterion and the strongest one (in which the rotational energy actually declines) is illustrated in Figure 1. Although a genuine evaluation of how large this effect might be requires numerical MHD simulations, simplified (simplistic?) analytic estimates indicate that even this criterion might be met when \( a_\ast \gtrsim 0.6 \) \[9\].
FIGURE 1. Fig. 1. The efficiency $\eta$ above which the rotational energy is diminished is shown by a solid curve. For comparison, the dotted curve shows the Novikov-Thorne efficiency; the dashed curve shows the efficiency at which $a_*$ does not change [1]. Both the solid and dashed curves are specific to the model of this section.
Because the spindown luminosity in this mechanism is intimately tied to accretion, its natural luminosity scale is the accretion luminosity scale, $\sim \dot{M}_c c^2$. However, effects such as mass clumping or magnetic reconnection could diminish it, possibly by sizable factors.

**VI** **FIELD LOOPS ANCHORED IN THE ERGOSPHERE, CLOSED INSIDE THE MARGINALLY STABLE ORBIT, AND NOT FORCE-FREE**

If, like the magnetic field on the horizon, the field strength in the plunging region is comparable to its value in the disk, it could in principle drive a substantial magneto-centrifugal wind. Unlike most MHD wind models, this one would be very far from time-steady, for the character of the field in this region changes in a free-fall time. In a time-steady wind, field lines must extend to infinity, for fluid elements have been travelling outward, carrying their magnetic flux with them, for effectively forever. By contrast, when the wind is highly variable, there is no requirement for any continuity in field structure between fluid elements expelled long ago and those just entering the wind today.

Nonetheless, the same dynamics that drive time-steady winds should also apply in this context. In fact, circumstances might be more propitious toward wind creation in this region than in the disk, for the Alfvén speed should be rather higher due to the similar strength magnetic field and the much smaller inertia (cf. the simulations reported in [20]). Local fluctuations in the angular momentum of plunging fluid elements could lead to shocks in this region, and strong local heating. Just as for the Blandford-Znajek mechanism, the characteristic luminosity scale is $\sim |\vec{B}|^2 r_g^2 c$, but the coefficient could be considerably less than unity.

**VIII ** **DISK WINDS**

Finally, it is useful as a standard of comparison to contrast these variations on the Blandford-Znajek mechanism with magneto-centrifugal winds launched from an accretion disk. Indeed, Livio et al. [19] claimed that the black hole spin energy reservoir would *always* be relatively unimportant because electromagnetic energy loss from the accretion disk would always exceed that from the black hole itself through the Blandford-Znajek mechanism. Their argument (an elaboration of one in [7]) hinged on a simple order-of-magnitude comparison: They estimated that the luminosity of a disk-driven wind would be generically $\sim O(10) c r_g^2 B_{pd} B_{td}$, where $B_{(p,t)d}$ are the poloidal and toroidal magnetic field components in the disk and the factor $O(10)$ accounts for the fact that the surface area of the disk is rather larger than the surface area of the black hole. Supposing that all currents supporting
magnetic fields are found in the disk, they then estimated $B_h \lesssim B_{pd}$, so that the disk luminosity would always be greater than the Blandford-Znajek luminosity.

However, this luminosity estimate is based on the implicit assumption that the $r-\phi$ component of the Maxwell stress in the disk wind has a vertical scale height comparable to the disk radius, not the (gas density-determined) disk thickness. This is a very strong assumption. At the moment, there is no way to rigorously evaluate its quality, but there are a few suggestive results in hand already from simulations [22,13]. Both of these indicate that the magnetic stresses decline vertically more slowly than does the gas density, but the scale height ratio is only a factor of two or three. If later work confirms these tentative results, the estimated magnetic wind luminosity of the disk would be greatly reduced. Mechanisms that derive their energy from the spin of the black hole, but do not put the energy into winds, would then rise in relative importance.

IX CONCLUSIONS

Virtually every idea anyone has suggested about how to extract spin energy from black holes involves magnetic fields anchored in the black hole’s ergosphere in one way or another. At the same time, the variations possible among these ideas (whether the field lines close at infinity or in the disk, whether they are anchored in the horizon or farther away, . . .) raise different questions and different opportunities. In discussions of these ideas, it is therefore important to distinguish carefully between them, for questions and criticisms raised about one version do not necessarily apply to the others. It has been the aim of this short paper to provide a convenient framework for making these distinctions, and to assess (in a necessarily subjective fashion) their current standing.

In order for further progress to be made, it is apparent that (at least) three key questions must be answered:
1.) Is it possible for energy and angular momentum to be carried away from a black hole when the field lines are anchored only in the horizon?
2.) Do black holes accumulate net magnetic flux (so that field loops close far away from the black hole and the structure changes on very long timescales) or does enough reconnection accompany accretion that field loops all close nearby, and the structure changes on the infall timescale?
3.) Can one think of the field as being force-free everywhere but in a small load, or is the plasma inertia qualitatively important?

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