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articles

Possible power source of Seyfert galaxies and QSOs

J. G. Hills

Department of Astronomy, University of Michigan, Ann Arbor, Michigan 48104

The possible presence of massive black holes in the nuclei of galaxies has been suggested many times. In addition, there is considerable observational evidence for high stellar densities in these nuclei. I show that the tidal breakup of stars passing within the Roche limit of a black hole initiates a chain of events that may explain many of the observed principal characteristics of QSOs and the nuclei of Seyfert galaxies.

STARS of mean density ρ are broken apart by the tidal force of a black hole of mass M if they pass within a distance

$$R_T \simeq (6M/\pi\rho)^{1/3} \quad (1)$$

of it. They are physically absorbed by the black hole if they pass it within the gravitational radius,

$$R_G = 2GM/c^2 = 2.95 \text{ km } (M/M_\odot) \quad (2)$$

R_T and R_T/R_G as functions of M are given in Table 1. I have assumed solar type stars in evaluating ρ in equation (1). For $M < 3\times10^8 M_\odot$, the tidal radius is greater than the gravitational radius so that stars are far more likely to be broken apart by the black hole than absorbed whole. By the time the black hole has grown to the critical mass $M_c \simeq 3\times10^8 M_\odot$, it may have produced several times its own mass in gas, by breaking up stars. This may be the source of the net outflow of gas observed in the nuclei of many galaxies.

Capture of gas by the black hole

Much of the gas produced by tidal breakup is captured and subsequently accreted by the black hole. The tidal breakup of a star requires energy which is ultimately supplied by a reduction in the kinetic energy of the star with respect to the black hole. The energy required to break up a star of mass m and radius R is

$$E_b = \tfrac{3}{4}(Gm^2/R) \quad (3)$$

for a polytrope of index $n=3$. The gas produced by the breakup of a solar type star is captured into a bound orbit by the black hole if the velocity of the star with respect to the black hole before their encounter was less than 535 km s^{-1}. Thus if the velocity dispersion of the stars in the nucleus of a galaxy is less than about 500 km s^{-1}, much of the gas produced by tidal breakup is captured and ultimately accreted by the black hole.

For the larger black holes, which have $R_T \sim R_G$, emission of gravitational radiation may also be important in reducing the kinetic energy of the star with respect to the black hole[1], but I shall not consider this additional loss mechanism here.

The gas produced by the breakup of a particular star may travel in a bound orbit to a large distance from the black hole, but eventually it must return to the original periastron distance of its stellar precursor. The mean semi-major of the gas produced by the breakup of a solar type star by a black hole of mass M is

$$\begin{aligned} r &= (2M/3m)R[1-2\langle V^2\rangle R/3Gm]^{-1} \\ &= 1.5 \text{ pc } (M/10^8 M_\odot)[1-\langle V^2\rangle/(535 \text{ km s}^{-1})^2]^{-1} \end{aligned} \quad (4)$$

where $\langle V^2\rangle$ is the mean-squared velocity of the stars in the galactic nucleus before their encounter with the black hole. Thus a $10^6 M_\odot$ black hole is surrounded by a massive gas cloud with a typical radius of 10^{-2} pc. The dust in this cloud strongly absorbs the ultraviolet radiation emitted by the gas being accreted by the blackhole. This may be responsible for the infrared emission of QSOs and Seyferts. The ionised gas in this large cloud may in turn be responsible for the forbidden-line radiation from QSOs.

Collisions among the gas clouds originating from the breakup of individual stars will release enormous amounts of energy. High velocity collisions between gas clouds can explain several phenomena associated with QSOs (ref. 2). The dissipation of energy from collisions reduces the effective semi-major axis of the gas until it is well within the Roche limit. Ultimately the gas must be accreted by the black hole unless it is stopped by rotation. Even if this occurs, the gyration radius of the gas must be much less than the Roche limit of the black hole unless the gas acquires additional angular momentum after the breakup of its stellar progenitors by interacting with the stars around the black hole. If the angular momentum is sufficiently large that the gas settles into a disk spinning around the black hole, viscous forces eventually will allow most of this gas to be accreted by the black hole as in an X-ray binary[3,4]. The accretion of gas by a massive black hole produces a large amount of ultraviolet radiation[4] which will tend to ionise the massive gas cloud around the black hole to produce the observed forbidden-line radiation and the cores of the permitted lines. The Doppler broadening

Table 1 Breakup and accretion of stars by a black hole

$M/M_\odot$	$R_T/R_\odot$	R_T/R_G	$L_E/L_\odot$ (Eddington Limit)	t_E (yr)	L/L_E ($\rho_s = 10^6 M_\odot$ pc^{-3})	t_D (yr)
10^1	4.31	1.02×10^5	3.2×10^5	1.6×10^9	1.5×10^{-3}	1.8×10^{11}
10^2	9.28	2.19×10^4	3.2×10^6	1.4×10^9	3.2×10^{-3}	8.4×10^{10}
10^3	20.0	4.72×10^3	3.2×10^7	1.2×10^9	7.0×10^{-3}	3.9×10^{10}
10^4	43.1	1.02×10^3	3.2×10^8	9.4×10^8	1.5×10^{-2}	1.8×10^{10}
10^5	92.8	2.19×10^2	3.2×10^9	7.3×10^8	3.2×10^{-2}	7.9×10^9
10^6	200.0	4.72×10^1	3.2×10^{10}	5.2×10^8	7.0×10^{-2}	3.3×10^9
10^7	431.0	1.02×10^1	3.2×10^{11}	3.2×10^8	1.5×10^{-1}	1.2×10^9
10^8	928.0	2.19	3.2×10^{12}	1.1×10^8	3.2×10^{-1}	2.7×10^8
3.2×10^8	1,374.0	1.00	1.0×10^{13}	0	4.8×10^{-1}	0

resulting from the rotation of the gaseous disk around the black hole and the large random velocities of dense colliding gas clouds near the black hole may be responsible for the large widths of the permitted emission lines in Seyferts and QSOs.

Luminosity and total energy

It is evident that most of the mass acquired by a black hole less massive than $M = 3\times10^8 M_\odot$ is from the accretion of gas rather than of whole stars. In this case the total energy released by this accretion is $E = 0.057–0.43Mc^2$ depending on the angular momentum of the black hole[5]. If $E = 0.2Mc^2$, the total energy released by the black hole by the time its mass reaches the critical value $M_c = 3\times10^8 M_\odot$ is $E_c = 1.1\times10^{62}$ erg $= 9\times10^{20} L_\odot$ yr. The maximum luminosity produced in the accretion of material by the black hole is given by the Eddington[6] limit,

$$L_E = 3.2\times10^4 L_\odot (M/M_\odot) \quad (5)$$

Thus the maximum possible luminosity produced by the accretion of gas by the black hole is $L = 10^{13} L_\odot$ which occurs when $M = M_c$. This is very near the peak observed luminosity of QSOs. The total energy release could power a quasar at this luminosity for as long as 9×10^7 yr.

If $E = 0.2Mc^2$, the rate of mass accretion by a black hole growing at the Eddington limit is

$$\mathrm{d}M/\mathrm{d}t_E = 1.1\times10^{-8} M_\odot \,\mathrm{yr}^{-1} (M/M_\odot) \quad (6)$$

In this case, the time required for a black hole to grow from a mass M to $M_c = 3\times10^8 M_\odot$ is

$$t_E = 9.3\times10^7 \ln(M_c/M)\,\mathrm{yr} \quad (7)$$

Table 1 shows t_E and L_E as a function of M. I assume that the luminosity of a QSO drops by orders of magnitude when its mass reaches $M = M_c$ at which point the stars are swallowed whole by the black hole rather than being broken apart tidally first.

I note (Table 1) that a seed black hole growing at the Eddington rate takes about 1.6×10^9 yr to increase its mass from about $10M_\odot$, a mass one may expect from stellar evolution, to $M_c = 3\times10^8 M_\odot$. So the observed number of QSOs in the Universe would decrease very rapidly after the Universe was more than 1.6×10^9 yr old. If the Universe is presently 1.2×10^{10} yr old, this critical time corresponds to $Z = 2.8$ in an Einstein–de Sitter universe. This is consistent with the observations of Schmidt[7] that the density of QSOs in the Universe increases as $(1+Z)^6$ up to $Z \simeq 2.5$ and then it either remains constant or decreases for larger Z.

Growth governed by stellar density

If the density of stars in the nucleus of a galaxy is not sufficiently large, this density rather than the Eddington limit governs the rate of growth and luminosity of the black hole. Assuming Newtonian mechanics, the rate at which gas is being produced by the breakup of stars passing within the Roche radius R_T of a black hole of mass M is

$$\frac{\mathrm{d}M}{\mathrm{d}t} = \rho_s \langle \sigma V \rangle$$

$$= 8.6\times10^{-16} \frac{M_\odot}{\mathrm{yr}} \left[\frac{\rho_s}{M_\odot/\mathrm{pc}^3}\right]\left[\frac{R_T}{R_\odot}\right]\left[\frac{M}{M_\odot}\right]\left[\frac{\mathrm{km\ s^{-1}}}{\langle V^2\rangle^{1/2}}\right]$$

$$= 1.9\times10^{-15} \frac{M_\odot}{\mathrm{yr}} \left[\frac{\rho_s}{M_\odot/\mathrm{pc}^3}\right]\left[\frac{\mathrm{g\ cm^{-3}}}{\rho}\right]^{1/3}\left[\frac{M}{M_\odot}\right]^{4/3} \times \left[\frac{\mathrm{km\ s^{-1}}}{\langle V^2\rangle^{1/2}}\right] \quad (8)$$

Here ρ_s is the stellar mass density in the vicinity of the black hole. It is assumed that at very large distances from the black hole the stars have a Maxwellian distribution of velocities with an r.m.s. velocity $\langle V^2\rangle^{1/2}$. In this derivation I have assumed that the mass and radius of a typical star broken apart by the black hole are much smaller than the mass and tidal radius of the black hole. I have also made use of the fact that the escape velocity at a distance R_T from the black hole is very much greater than $\langle V^2\rangle^{1/2}$.

If all the gas produced in the breakup of the stars is accreted by the black hole integrating equation (8) suggests that the time required for a black hole to grow from a mass M to a mass M_c is

$$t_D = 1.7\times10^{15}\,\mathrm{yr} \left[\frac{M_\odot\ \mathrm{pc}^{-3}}{\rho_s}\right]\left[\frac{M_\odot}{M}\right]^{1/3}\left[1-\left(\frac{M}{M_c}\right)^{1/3}\right] \times \left[\frac{\langle V^2\rangle^{1/2}}{\mathrm{km\ s^{-1}}}\right] \quad (9)$$

Here I have assumed solar type stars so that $\rho = 1.41$ g cm^{-3}. Table 1 shows t_D as a function of M for a galactic nucleus in which $\rho_s = 10^6 M_\odot$ pc^{-3} and $\langle V^2\rangle^{1/2} = 225$ km s^{-1}, the observed velocity dispersion in the nucleus of M31 (ref. 7). The table also gives the luminosity, L, in units of the Eddington luminosity, L_E. Here I have used $L = 0.2\,(\mathrm{d}M/\mathrm{d}t)\,c^2$. I note that at a fixed M, $(t_D)^{-1} \propto L \propto \rho_s/\langle V^2\rangle^{1/2}$.

Thus simple scaling gives L and t_D for a galactic nucleus with any other value of ρ_s and $\langle V^2\rangle^{1/2}$. It is not likely that $\langle V^2\rangle^{1/2}$ differs very much from one nucleus to another, but there are sizable differences in ρ_s. Assuming $M/L = 10M_\odot/L_\odot$, the data of Kinman[8] indicate that $\rho_s \simeq 10^5 M_\odot$ pc^{-3} within 1 pc of the nucleus of M31. For the nucleus of our Galaxy infrared observations[9] indicate that $\rho_s \simeq 7\times10^6 M_\odot$ pc^{-3} within 0.5 pc of its centre. Schwarzschild[10] finds indirect evidence of $\rho_s = 2.2\times10^8 M_\odot$ pc^{-3} within 3.5 pc of the nucleus of the Seyfert galaxy NGC4151. For a black hole to have had enough time to grow to $M_c = 3\times10^8 M_\odot$ within the age of the Universe (which I assume to be 1.2×10^{10} yr), the initial mass of the black hole had to exceed $M_i \simeq 10^7 M_\odot$ in the nucleus of M31, $M_i \sim 10^2 M_\odot$ in the nucleus of our Galaxy, and $M_i \lesssim 1 M_\odot$ in the nucleus of NGC4151. The very low value of M_i for NGC4151 suggests that Schwarzschild's estimate for ρ_s may be somewhat large. From the values of L/L_E in Table 1, a black hole in the nucleus of M31 could never break up stars fast enough to allow it to radiate at the Eddington limit. In our Galaxy a black hole would radiate at the Eddington limit when its mass is equal or greater than $M_E = 10^7 M_\odot$; for NGC4151 this occurs when $M_E = 10^2 M_\odot$. When $M > M_E$ the breakup of stars by the black hole produces gas faster than the Eddington limit allows it to be accreted by the black hole. In these cases, we may expect the surplus gas to be driven out of the nucleus by radiation pressure. Such outpouring of gas is observed in several Seyferts. This mass loss can be prodigious. If $\rho_s = 10^7 M_\odot$ pc^{-3}, the mass loss rate approaches 12.5 $M_\odot$ yr^{-1} as M approaches M_c. The mass loss rate as well as the luminosity would then drop effectively to zero for $M > M_c$.

If present ideas on the evolution of massive upper main sequence stars are correct, we may reasonably expect black holes with masses of about $M_i = 10 M_\odot$ to form in almost every galactic nucleus. Such black holes would have had sufficient time in our Universe to grow to $M_c = 3\times10^8 M_\odot$ if the stellar density in the galactic nucleus exceeds $(\rho_s)_0 = 1.5\times10^7 M_\odot$ pc^{-3} (Table 1). This watershed density is very sharply defined. If ρ_s is even 2–3 times less dense than this, the minimum mass of a seed black hole which has sufficient time to grow to M_c in the age of the Universe increases sharply to 10^2–$10^3 M_\odot$ which is much greater than the masses of black holes formed by normal stellar evolution. If $\rho_s > 2$–$3\ (\rho_s)_0$, almost any black hole has time to grow to M_c. Thus any QSO or Seyfert nuclei present today are likely to have a stellar density of about $1.5\times10^7 M_\odot$ pc^{-3} or slightly greater. At this density the black hole radiates at the Eddington limit for $M > M_E = 10^5 M_\odot$. If $\rho_s \geq 10^8$–$10^9 M_\odot$ pc^{-3} any black hole with $M > 10 M_\odot$ would grow at the rate governed by the Eddington limit. These black holes would have reached $M = M_c$ at about $Z = 2.8$ in the observable Universe. Thus there should be a maximum in QSO activity at this epoch followed by a rapid decline for lower values of Z. This is consistent with observations[7].

Two refinements may increase somewhat the rate at which stars are swallowed by the black hole. One complication is the possible formation of a high density stellar cusp around a black hole embedded in a dense stellar nucleus[11]. Another is the shrinking of the entire stellar nucleus as its stars are captured by the black hole. How this arises can be seen by considering the effect of these captures on a typical star in the nucleus. Some of the stars captured by the black hole are in orbits in which, on the average, they are further from the black hole than the test star. At some point in time each of these stars crosses inside the orbit of the test star never to return. Thus, as these stars are accreted by the black hole, the test star and the other remaining stars in the nucleus feel a progressively increasing gravitational potential. To satisfy the virial theorem, the velocity dispersion of the stars in the nucleus must increase. This added kinetic energy is supplied by a net shrinking of the effective radius, R_N, of the galactic nucleus. The increase in the velocity dispersion of the stars in the nucleus would by itself decrease the rate of accretion by the black hole, but this is more than compensated by the increased stellar density in the nucleus. I note that $\langle V^2\rangle \propto R_N^{-1}$ while $\rho_s \propto R_N^{-3}$ (actually, $\rho_s \simeq R_N^{-2}$ when stars accreted by the black hole are considered). From equation (8), this results in a net increase in the rate of accretion by the black hole compared to the case where we ignore both the increased density and the increased velocity dispersion in the galactic nucleus. This effect can only become important as the mass of the black hole approaches that of the galactic nucleus or $M \sim 10^8 M_\odot$.

Orbit diffusion and relaxation time

In the previous section, particularly in the derivation of equation (8), I made the implicit assumption that the rate of orbit diffusion resulting from the random gravitational encounters between stars is sufficiently fast that at any given time the number of stars with orbits that pass within the Roche radius, R_T of the black hole is nearly the same as it would be if there was no breakup of these stars. This is true as long as the total mass M_n of the stars in the galactic nucleus is large compared to the mass M of the black hole. In the absence of orbit diffusion, the velocity distribution of the stars at a given point in the nucleus would be isotropic (if the stellar distribution is relaxed) except for those velocity directions corresponding to orbits which cross R_T. On average, the fraction, F_v, of the velocity directions occupied by these orbits is the fraction of all stars in the nucleus that have orbits which intersect R_T. Thus as long as $F_v \lesssim 1$, a very small amount of velocity randomisation is sufficient to fill in these small holes in the otherwise isotropic velocity distribution and consequently to repopulate the orbits that cross R_T. In other words the time necessary to repopulate these orbits is very small compared to the relaxation time[12],

$$t_R = V^3/4\pi G^2 m^2 n \ln N \qquad (10)$$

which is the time necessary to completely rerandomise the velocities. In the typical galactic nucleus of interest, $V = 225$ km s^{-1}, $n = 10^7$ stars pc^{-3}, $m = 1 M_\odot$, $M_n = 10^8 M_\odot$, and $N = M_n/m = 10^8$. In this case $t_R = 3\times10^8$ yr.

The rate of orbit diffusion is sufficiently large to result in there being no significant depletion of the stars with orbits that cross R_T as long as the characteristic time for the consumption of all the stars in the galactic nucleus by the black hole is greater than the relaxation time; that is, the diffusion rate is sufficient as long as $M_n/|dM_n/dt| > t_r$. If $M_n \sim 10^8 M_\odot$, this condition is satisfied until the mass of the black hole has grown to $M \sim M_c$. Thus the assumption of an adequate rate of orbit diffusion implicit in the previous section is valid.

A massive black hole at the galactic centre?

Any such object would have to be very modest to produce no more than the observed nonstellar radiation emitted from the nucleus. If gas accretion by this object produces all the ionising radiation required to maintain the H II region at the galactic centre, the luminosity of this radiation is $L_i = 1.9\times10^8 L_\odot$ (ref. 13) if the energy per photon is 2 rydberg. From Table 1, with $\rho_s = 7\times10^6\ M_\odot$ pc^{-3}, this requires a black hole mass $M \simeq 4\times10^4\ M_\odot$. Equation (3) requires that the radius of the cloud of gas and dust produced by the breakup of stars by a black hole of this mass be $\langle r\rangle \simeq 10^{-3}$ pc. This is at least consistent with the observed "point" infrared source at the galactic centre which has a radius of less than 0.1 pc and a luminosity of $3\times10^5\ L_\odot$. (ref. 8) $M \simeq 4\times10^4\ M_\odot$ is likely to be a firm upper limit on M. The ionisation could be due to something else. If M is this large, it would only take another 1.5×10^9 yr to grow to M_c. It is unlikely (but of course not impossible) that we would catch the black hole just at this crucial stage in its growth. A more troublesome point is the fact that the initial mass of this object had to be about $10^2 M_\odot$ to have reached $M = 4\times10^4$ by the present epoch. This initial mass seems rather large from the standpoint of stellar evolution. If the luminosity of the infrared source results from the accretion of stars broken up by the central black hole, its mass must be at least $M \simeq 3\times10^2 M_\odot$. This is probably the minimum mass of the central black hole. The initial mass of this black hole need only have been slightly greater than $10 M_\odot$ to have allowed it to grow to $3\times10^2 M_\odot$ in 1.2×10^{10} yr. It will take another 10^{10} yr for such a black hole to grow to $M = M_c$.

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Dredged basalts from the mid-oceanic ridge north of Iceland

Finn Dittmer
Ocean Study Group, Roskilde Universitetscenter, 4000 Roskilde, Denmark

Stanley Fine, Mogens Rasmussen, John C. Bailey & John Campsie
Ocean Study Group, Universitetets, Mineralogisk-Geologiske Instituter, og Mineralogsik Museum, Øster Voldgade 5–7, 1350 Copenhagen K, Denmark

An examination of the concentrations of the incompatible elements in tholeiites from the Kolbeinsey Ridge indicates that plume magmas have probably not overflowed into that area from Iceland or Jan Mayen Island. The results also show that the chemistry of plume magmas may be intimately related to regional geochemical patterns, and it is possible that Iceland represents a 'weak spot', or decompression focus, which causes complex flow patterns within and from the asthenosphere.

We present here petrological data on dredge samples recovered from the mid-oceanic Kolbeinsey Ridge north of Iceland. Most of the materials described were recovered during two cruises in 1972–73 in the North Atlantic.

Iceland and Jan Mayen Island represent subaerial expressions of recent volcanism associated with the Iceland Plateau and Kolbeinsey Ridge. Iceland consists of an extensive, largely tholeiitic, flood-basalt plateau built up in about the last 30 Myr and cut by three neovolcanic zones containing tholeiitic, alkalic and other basalt types[1-3]. Its great volumes of lava, their high discharge rates and relatively high contents of K, Ti, P, rare earth elements (REE), $^{87}Sr/^{86}Sr$ and so on have been considered expressions of a hot, rising mantle plume[2,4-6].

The older Faeroes and East Greenland basalts also have K, P, and Ti values comparable to those found in Iceland[7], and this may be[8] the result of plume activity in the early Tertiary.

Jan Mayen Island consists of ankaramites, alkali basalts and their differentiates together with their explosive equivalents; tholeiitic flows are unknown. Subaerial material only dates back to 0.5 Myr BP but construction of the island's base could be assigned to the mid-Tertiary or early Pleistocene[9,10]. A plume origin has also been proposed for Jan Mayen Island[11]. The two islands lie at the limits of the Iceland Plateau, a generally smooth (presumably basaltic) platform at a depth of less than 1 km (refs 12 and 13). Transecting this plateau is the Kolbeinsey Ridge, an accretion centre erupting basalts at its axis and with minimal sedimentary cover[13-15]. Possible extensions of recent Icelandic and Jan Mayen plume activity may thus be sought along the Kolbeinsey Ridge.

The Kolbeinsey Ridge (66° 30′N–71° 50′N) extends, with a generally NNE trend, for a distance of approximately 550 km from Iceland to the Jan Mayen Fracture Zone and cuts the Icelandic Plateau at depths of less than 1,300 m (Fig. 1). The ridge begins as a subdued N–S trending rift valley and corresponding magnetic anomaly which gradually develops on the outer northern Icelandic Shelf. The shallow ridge axis surfaces at the islet of Kolbeinsey, 67° 08′N (ref. 16). North of Kolbeinsey Islet as far as the Spar Fracture Zone (69°N) the ridge strikes N 28°E and has a step-like structure lacking a central rift valley, it is similar in morphology to the northern Reykjanes Ridge. In addition, this ridge segment also resembles the Reykjanes Ridge in spreading rate, free-air and Bouguer anomalies, and magnetic anomalies[14]. Remanent magnetisation values on drilled cores from pillow basalts from this ridge segment obtained by our shipboard spinner magnetometer averaged 0.05 e.m.u., similar to the average values found for pillow basalts on the Reykjanes Ridge[17]. Furthermore, thin sections show phenocryst assemblages of plagioclase–olivine–clinopyroxene (dredges 7–72, 9–72, see Table 1) which are also characteristic for pillow basalts from the Reykjanes Ridge[4,18]. Thus, the conclusion[14] that this segment of the Kolbeinsey Ridge seems to be the northern counterpart of the Reykjanes Ridge is probably correct.

North of the Spar Fracture Zone, the Kolbeinsey Ridge is offset about 30 km to the east and then takes on the classical mid-oceanic ridge structure with a well developed rift valley trending N 14°E. An unnamed fracture zone at 70° 52′N delimits the southern side of Eggvin Bank, a broad seamount transected by the ridge axis. From the top of Eggvin Bank, only 50 m below the waves, the ridge gradually deepens towards the Jan Mayen Fracture Zone which reaches depths greater than 2 km.

Chemistry of Kolbeinsey Ridge tholeiites

Table 2 presents the concentrations of K_2O, TiO_2, P_2O_5 and Sr in tholeiites from the Kolbeinsey Ridge; these have been plotted against the morphology of the ridge axis (Fig. 1). Tholeiites from the ridge axis between 67° and 70° 35′N have similar low contents of the incompatible elements; low levels of K_2O, TiO_2, P_2O_5 and Sr characterise fresh mid-oceanic ridge tholeiites[19-21]. In detail, the levels of K_2O, P_2O_5, Sr, and to a lesser extent TiO_2, are generally lower than oceanic ridge tholeiites south of the Charlie Gibbs Fracture Zone[20], and are more comparable to contents of the central and southern Reykjanes Ridge[4,18,22].

Eggvin Bank also consists of tholeiitic basalts but their elemental concentrations deviate sharply from the 'background level' determined for the ridge to the south. Figure 1 shows that the concentrations of K_2O, TiO_2, P_2O_5 and Sr rise to a maximum at the upper parts of Eggvin Bank. This peak in no way corresponds to the presence of highly evolved or fractionated basalts: the FeO*/MgO ratio of Eggvin Bank tholeiites is comparable with the ratio for tholeiites from the adjacent ocean floor to the south. Moreover, fractional crystallisation would lead to a similar passive concentration of elements such as K and P which enter early mineral phases in negligible amounts. In fact, K_2O concentrations are increased by a factor of seven over ridge values whereas P_2O_5 is only increased twofold. The concentration of alkali and volatile elements within a high-level magma chamber[23] of Eggvin Bank is a possibility. Alternatively, the lower TiO_2/P_2O_5 ratio of these summit tholeiites