A Link Between Cosmic Ray Flux and the Formation of Star Clusters During Galaxy-Galaxy Interactions

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ABSTRACT

If the magnetic fields that are responsible for supporting molecular clouds against gravitational collapse in our Galaxy were to instantaneously decouple from the gas (and dust), the clouds probably would undergo a burst of star formation. In many cases, the efficiency of star formation would be sufficiently high to produce a gravitationally bound cluster of stars and, in so doing, would generate a luminosity per unit mass that rivals any starburst phenomenon. We propose that this is precisely what happens during galaxy-galaxy interactions as gravitationally bound molecular gas clouds enter an environment in which the cosmic ray flux – and resulting fractional ionization of the dense cloud gas – drops to a level that is much lower than it was in the undisturbed galaxy disk. We propose that the decoupling of magnetic fields from giant molecular clouds that results from this drop in the fractional ionization is the mechanism that is primarily responsible for the formation of massive, bound star clusters and, in turn, the starburst phenomenon during strong galaxy-galaxy interactions. Because it does not rely upon violent gas-dynamical triggers, this mechanism may also explain how bound star clusters are formed in systems like the Magellanic Clouds during relatively mild tidal interactions with a larger galaxy companion.

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1. Introduction

High resolution images taken with the wide-field camera on the HST have made it abundantly clear that galaxy-galaxy interactions provide an environment that is conducive to the formation of massive, bound clusters of stars (Meurer *et al.* 1995; Schweizer *et al.* 1996; Miller *et al.* 1997; Johnson *et al.* 1999; Zepf *et al.* 1999; Whitmore *et al.* 1999). (See Whitmore 2000 for a complete list of interacting galaxies with young star clusters.) Apart from their very young ages, the observed properties of these clusters, both individually and as groups, resemble in many respects the properties of globular clusters. This has led to speculation that the conditions under which globular cluster systems formed in galaxies when the universe was relatively young were very similar to the conditions that we can now study in merging systems at the present epoch (Lee, Schramm & Mathews 1995; Elmegreen & Efremov 1997). Indeed, before the first HST images were collected, Ashman & Zepf (1992) and Zepf & Ashman (1993) predicted that newly formed globular clusters would be found in the cores of merging galaxy systems.

Here we are not interested in debating to what degree the newly formed star clusters in merging galaxy systems actually resemble globular clusters, or to what degree the conditions under which star clusters formed in the early universe resemble the conditions one finds in merging systems at the present epoch. We wish, instead, to focus on the more generic, but firmly established realization that galaxy-galaxy interactions provide at least one environment that is conducive to the formation of massive, bound star clusters (Meurer *et al.* 1995; Schweizer & Seitzer 1998). This is in stark contrast to the normal conditions in undisturbed gas-rich galaxy disks where star formation is a relatively inefficient process and, in particular, where it is apparently rather difficult to make bound star clusters of significant mass (Roberts 1957; Blaauw 1964; Miller & Scalo 1979; Brush & Sanders 1983; Lada, Margulis & Dearborn 1984; Battinelli *et al.* 1994). Following the arguments of Lada, Margulis & Dearborn (1984), there must be some process associated with galaxy interactions that raises the star formation efficiency from its typical value of only a few percent or less in the disk to $\gtrsim 50\%$. What is this physical process?

It has been suggested that gas-dynamical "triggers" such as cloud-cloud collisions (Scoville, Sanders, & Clemens 1986; Olson & Kwan 1990; Habe & Ohta 1992; Lee, Schramm & Mathews 1995) or environmental over-pressures (Jog & Solomon 1992; Elmegreen & Efremov 1997) are primarily responsible for the high efficiency of star formation and/or the formation of bound star clusters in interacting systems. We suggest, instead, that bound clusters form because there is a relatively rapid decoupling of the magnetic field from molecular clouds as they arrive at the scene of an interaction as part of a gas-rich host. This idea seems reasonable because: (1) The gravitationally bound molecular clouds that are found in the disk of our Galaxy (as well as in most other gas-rich disk galaxies) provide ideal seeds for the formation of star clusters having the range of masses that are presently seen in merging galaxy systems; (2) it is the coupling of the magnetic field to the partially ionized gas that limits the rate at which stars are presently forming in molecular clouds in the disk of our Galaxy; and (3) models of cosmic-ray propagation in our Galaxy predict that the cosmic-ray flux — which controls the ionization rate in dense molecular clouds — drops significantly outside of the disk of the Galaxy. Hence, if an existing molecular cloud were displaced from the galaxy disk to a region where the cosmic-ray flux is sufficiently low or, more likely, if the relatively ordered, large-scale magnetic field that is responsible for confining cosmic rays to the galaxy becomes sufficiently disorganized so that the cosmic rays no longer remained confined near the disk, the star formation rate throughout the cloud would increase significantly and the cloud would spontaneously transform into a bound cluster of stars. This mechanism for raising the efficiency of star formation in molecular clouds does not rely upon violent gas-dynamical processes which, in reality, tend to disrupt clouds and thereby actually lower the overall efficiency of star formation (Elmegreen 1998). (In an appendix, we explain why isothermal cloud collisions, in particular, cannot be relied upon to make stars with high efficiency.) In what follows, we address in more detail each logical step of this scenario.

2. The Properties of Molecular Clouds

Molecular clouds contain enough mass and are of about the right size to produce clusters of stars whose overall properties resemble the populations of star clusters with which we are familiar, up to and perhaps rivaling the largest globular clusters (Jog & Solomon 1992; Harris & Pudritz 1994; Schweizer & Seitzer 1998). Over a wide range of masses ($10^3 - 10^6 M_\odot$), the clouds appear not only to be gravitationally bound structures, but structures in which the energy density in the gravitational field far surpasses the thermal energy density of the gas (Larson 1981; Sanders, Solomon, & Scoville 1984; Sanders, Scoville, & Solomon 1985). So, by a substantial margin the available thermal pressure is insufficient to support the clouds against gravitational collapse. In addition to this, molecular clouds exhibit a tremendous amount of internal structure. As Larson (1981) summarized two decades ago, at gas densities $1\,\mathrm{cm}^{-3} \lesssim n_{H_2} \lesssim 10^4\,\mathrm{cm}^{-3}$ ($3\times10^{-24}\,\mathrm{g\,cm}^{-1} \lesssim \rho \lesssim 3\times10^{-20}\,\mathrm{g\,cm}^{-3}$), nonlinear-amplitude density fluctuations are clearly visible on linear scales as small as 0.1 pc, with corresponding mass scales less than or on the order of the thermal Jeans mass $M_J \sim 1 M_\odot$. Naively, then, we should expect to be seeing molecular clouds transformed into large clusters of stars on a free-fall timescale,

$$t_{ff} \sim (G\rho)^{-1/2} \sim 10^6 \mathrm{yr} \left(\frac{\mathrm{n_{H_2}}}{4.5\times10^3 \mathrm{cm}^{-3}}\right)^{-1/2} . \qquad (1)$$

It is equally clear from a variety of different lines of argument, however, that molecular cloud structures in the disk of our Galaxy are not changing drastically on a free-fall timescale and that, overall, the process by which they make stars is very inefficient (Scoville, Sanders, & Clemens 1986; Tan 2000). At all scales larger than the thermal Jeans mass, molecular clouds exhibit internal (supersonic) velocity dispersions that are sufficiently large to support the clouds against free-fall collapse (Larson 1981). In the terms used by Jog & Solomon (1992), these large velocity dispersions give rise to an "effective temperature" that is significantly larger than the normal gas temperature and is sufficient to keep the clouds in virial balance. Although it has been argued that these velocities arise from hydrodynamical turbulence (Fleck 1983; Henriksen & Turner 1984), the growing consensus is that molecular clouds are supported against gravitational collapse by the interstellar magnetic field and that the observed velocities are associated with MHD waves which, while being supersonic, are generally sub-Alfvénic (Shu, Adams, & Lizano 1987; Myers 1987; Mouschovias 1987; Myers & Goodman 1988; Fleck 1988; Myers & Khersonsky 1995). It is quite reasonable to argue that the support comes from magnetic fields because, within molecular clouds, the energy density in the magnetic field is comparable to the energy density in the gravitational field and because the gas and dust in these clouds is partially ionized and therefore coupled to the field.

Although molecular clouds, as a whole, appear to be dynamically stable against gravitational collapse, star formation is nevertheless taking place in their densest cores. As outlined by Shu, Adams, & Lizano (1987) this happens because dense cloud cores also represent regions of relatively low fractional ionization, so the gas in these cores gradually loses magnetic flux via a process called ambipolar diffusion. As is discussed in more detail below, this process is generally slow relative to the local free-fall time. As a result, star formation proceeds in a relatively inefficient manner in molecular clouds. The scenario that we are proposing here to explain why bound clusters of stars are able to form during interactions between galaxies depends critically on this idea that, in a normal spiral galaxy disk, the interstellar magnetic field is responsible for supporting molecular clouds against dynamical collapse. (That is to say, if you are not convinced that the interstellar magnetic field is primarily responsible for providing this support, then you should automatically reject our cluster formation scenario as well.) We are suggesting that, during a galaxy-galaxy interaction, the magnetic field is able to decouple from much of the molecular gas and, as a result, gravitational fragmentation of molecular clouds is able to proceed on a free-fall time down to scales that are governed by the thermal Jeans mass (*i.e.,* $\lesssim 1M_{\odot}$). In an effort to ascertain under what conditions such a scenario might take place, we must review what conditions and physical processes dictate the ambipolar diffusion time in molecular clouds.

3. The Cause and Effects of Ionization

3.1. Ionization Due to Cosmic Rays

As has been discussed in considerable depth by McKee (1989; see especially Appendix B), ionization of the gas and dust in molecular clouds occurs primarily from two sources: high energy photons from the interstellar radiation field, and cosmic rays. In the outermost layers of any cloud complex, photoionization is dominant. But, as a direct result of this attenuation of the flux of ionizing photons, the process of ionization becomes dominated by cosmic rays in the inner regions of most clouds. McKee (1989) has estimated that the crossover typically occurs at an extinction of $A_V \sim 4$. The focus of our discussion will be on these inner cloud regions, which are generally of higher density and therefore more conducive to star formation in the first place. Henceforth, we will assume that cosmic rays are the primary source of ionization in molecular clouds.

In their pioneering paper, Spitzer & Tomasko (1968) examined to what degree cosmic rays are responsible for the ionization and heating of interstellar gas. Taking what was known about the measured spectrum of cosmic ray protons at the time, they derived an empirical expression for the local interstellar spectrum (LIS) of cosmic rays, $j(E)$ [expressed in units of particles $\mathrm{cm}^{-2}\mathrm{ster}^{-1}\mathrm{s}^{-1}(\mathrm{GeV/nucleon})^{-1}$], that peaks around 100 MeV/nucleon and drops off at high energies as $\mathrm{E}^{-2.6}$. They then calculated an ionization rate via the integral,

$$\zeta_{CR} = \frac{5}{3}\int 4\pi j(E)\sigma dE\,, \tag{2}$$

where σ is the cross-section for ionization given by Bethe (1933), obtaining a value of $\zeta_{CR} = 6.8\times10^{-18}\mathrm{s}^{-1}$. It is this determination — or, rather, its value rounded to the nearest power of ten, $10^{-17}\mathrm{s}^{-1}$ — that has been widely utilized over the past thirty years to estimate fractional ionization levels and ambipolar diffusion timescales (as discussed below) in dense molecular clouds (Dalgarno & Oppenheimer 1973; Spitzer 1978; Elmegreen 1979; Shu, Adams, & Lizano 1987; McKee 1989; Shu 1992; Mestel 1999).

Despite its widespread and frequent use, we are not often reminded that considerable uncertainty accompanied the approximate formula that Spitzer & Tomasko used for $j(E)$ in their calculation of ζ_{CR}. The astrophysics community's widespread acceptance of this value almost certainly comes from the realization that it is consistent with values of ζ that are required to explain measured values of the ionization fractions of various molecular species in dense clouds (Black *et al.* 1990), and that the total energy density in cosmic rays that one derives from Spitzer & Tomasko's $j(E)$ is approximately equal to the energy density in the interstellar magnetic field (McKee 1989).

More reliable determinations of $j(E)$ can be obtained from the multitude of cosmic-ray experiments that have been conducted over the past three decades (Webber 1998). But even today a determination of the LIS of cosmic-ray nuclei from measurements of cosmic rays in the solar system is uncertain, particularly at energies below ~ 100 MeV. This is because the solar system is significantly shielded from interstellar cosmic rays by the heliosphere. As measured from Earth environs, the cosmic ray flux is modulated on an 11 year cycle in response to sunspot activity (Sao *et al.* 1991; Webber 1998), and one must build a model of the interaction of cosmic rays with the heliosphere before a reasonable estimate of the LIS can be ascertained. Webber's (1998) most recent determination of $j(E)$ for cosmic-ray nucleons peaks at a lower energy and drops off with a steeper slope at high energies than the function used by Spitzer & Tomasko (1968). He also has derived $j(E)$ for electrons which, he points out, play a nonnegligible role in ionization. Using eq. (2) and including contributions from different particle species, he derives a value of $\zeta_{CR} = (3{-}4)\times10^{-17}\mathrm{s}^{-1}$. This is a factor of five to six larger than the value first quoted by Spitzer & Tomasko (1968), but still within an order of magnitude of the frequently quoted value of $10^{-17}\mathrm{s}^{-1}$. Hence, although we would argue that Webber's determination should be adopted by the community as an improvement over the result reported by Spitzer & Tomasko (1968), we will continue to use the canonical value because an order of magnitude estimate of $\zeta_{\rm CR}$ is sufficient in the context of our present discussion.

Over the past three decades we have gained a much better understanding of the origin and acceleration of cosmic rays, as well as a better understanding of the links between cosmic rays and photon radiation in our Galaxy. This, along with a much improved cosmic ray database, has made it possible to construct models of cosmic ray propagation that apply not only in the vicinity of the solar system but throughout the Galaxy (Webber & Rockstroh 1997; Strong & Moskalenko 1998). These models permit us now to draw conclusions about the cosmic ray flux — and the ionization rate determined from it — outside of our local region of the interstellar medium. We will return to this subject in §4, below.

3.2. Ambipolar Diffusion

In a partially ionized gas cloud, only the ionized particles have their motions directly influenced by the presence of a magnetic field. The neutral particles do not notice the field directly and want to respond, instead, to the local gravitational acceleration – *i.e.*, the neutral gas wants to flow toward any local minimum in the gravitational field and form a more compact object. But, as has been explained clearly by Shu (1992) (see also Mestel & Spitzer 1956, Nakano 1984, and Mestel 1999), the neutral particles are influenced by the magnetic field indirectly through collisions with the charged particles and these collisions prevent the neutral gas from simply free-falling in response to the gravitational field. Instead, the neutrals can only "drift" past the charged particles – and, correspondingly, the magnetic flux can only diffuse out of the gas cloud – at a rate and on an "ambipolar diffusion" timescale, t_{AD}, that is governed by this collision process.

Drawing on the works of Elmegreen (1979) and Shu (1983), as summarized by Shu, Adams, & Lizano (1987) (see also Dalgarno & Oppenheimer 1973, Spitzer 1978, Mouschovias 1987, and McKee *et al.* 1993), we find that for typical molecular cloud sizes, magnetic field strengths, and densities, t_{AD} is directly proportional to the ionization fraction of the gas, $x \equiv n_e/n_{H_2}$, with a proportionality constant such that,

$$t_{AD} \sim 8 \times 10^6 {\rm yr} \left(\frac{x}{1.5 \times 10^{-7}} \right) . \tag{3}$$

If the ionization process is dominated by cosmic rays, as we have argued is the case, the ionization fraction, in turn, will depend on the metallicity depletion δ, the cosmic-ray ionization rate $\zeta_{\rm CR}$, and the number density of molecules through the expression,

$$x \sim x_{\rm CR} \approx 1.5 \times 10^{-7} \left(\frac{\delta}{0.1} \right)^{1/2} \left(\frac{\zeta_{\rm CR}}{10^{-17}{\rm s}^{-1}} \right)^{1/2} \left(\frac{n_{H_2}}{4.5 \times 10^3 {\rm cm}^{-3}} \right)^{-1/2} . \tag{4}$$

Hence, independent of the molecular cloud density,[1] the ratio of the ambipolar diffusion timescale

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[1]Elmegreen (1979) has warned that, at sufficiently high number densities, the neutral particle drift rate is limited

to the free-fall time is,

$$\frac{t_{AD}}{t_{ff}} \sim 8 \left(\frac{\delta}{0.1}\right)^{1/2} \left(\frac{\zeta_{\rm CR}}{10^{-17}{\rm s}^{-1}}\right)^{1/2} , \tag{5}$$

and we see that, under typical conditions in the disk of our Galaxy ($\delta \sim 0.1$ and $\zeta_{\rm CR} \sim 10^{-17}{\rm s}^{-1}$), the ambipolar diffusion timescale is an order of magnitude larger than the free-fall time. [As we are reminded most forcefully by Myers & Khersonsky (1995), this ratio of timescales is even larger — and the star formation efficiency correspondingly lower — in the low density envelopes of molecular clouds where photoionization is operating as well.] This means that star formation can proceed only on a timescale that is significantly longer than the natural dynamical time of the gas cloud; *i.e.,* in one crossing time, only a relatively small fraction of the cloud mass will be converted into stars. As a result, the first stars that form from the molecular cloud will have orbital trajectories that are governed primarily by the remaining cloud mass and will, in no sense, act as a gravitationally bound star cluster. Then, statistically, before more than a few thousand stars have had a chance to form, one or more massive stars will not only form but will complete their entire life-cycle. The radiation, winds and, ultimately, explosion of the most massive stars will evaporate or disrupt the gas cloud, leaving behind an unbound "association" of stars (Lada, Margulis & Dearborn 1984; Lada & Lada 1990).

Relation (5) makes it immediately apparent, however, that if the cosmic ray ionization rate was set to a value that is two (or more) orders of magnitude less than its canonical value of $10^{-17}{\rm s}^{-1}$ — for example, closer to its intergalactic value $\zeta_{\rm IG} \sim 3 \times 10^{-21}{\rm s}^{-1}$ (see §4, below) — the magnetic field would decouple from the gas quickly enough to permit free-fall collapse to proceed. Then, at least in principle, a large fraction of the cloud mass could be converted into stars in a single crossing time. The resulting high efficiency of star formation would generate a luminosity per unit mass that rivals any starburst phenonemon (Heckman 1994) and a massive, bound cluster could remain behind after the residual cloud gas has been evaporated away.

4. Conditions Outside a Galaxy Disk

Because cosmic rays make a significant contribution to the reservoir of energy in the interstellar medium, a considerable amount of effort has been invested in the development of models of their origin and propagation. As discussed by Longair (1992; see especially chapter 20), propagation scenarios historically have been based on a so-called "leaky box" model, where the size and extent of the "box" has been more or less associated with the Galaxy's disk. Once produced, cosmic rays are trapped for a time by the Galaxy's large-scale magnetic field, but eventually they escape. The

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by collisions between the neutrals and dust via a process that prevents t_{AD} from dropping below some limiting value t_{limit}. For the typical molecular cloud conditions being considered here, $t_{limit} \sim 10^6$ yr, which is sufficiently small as to not conflict with the primary conclusions of this paper. Hence, this "dust limit" will not enter into our discussions further.

timescale on which a given nucleonic species leaks out of the Galaxy is a function of kinetic energy, but the most abundant and lowest energy particles characteristically remain trapped for $10^7 - 10^8$yr and particles with energies exceeding 100 GeV/nucleon generally are able to escape entrapment altogether.

It was pointed out long ago by Ginzburg, Khazan & Ptuskin (1980) that a more realistic model involves diffusive propagation from the cosmic-ray sources, and more recently it has been realized that multi-component, diffusion models are required in order to satisfactorily explain measured primary-to-primary and secondary-to-primary cosmic-ray abundance ratios. These models have been further refined to account for the detailed all-sky maps of γ-ray emission that have been provided over a wide range of photon energies by the instruments onboard the Compton Gamma-Ray Observatory (Hunter, Kinzer & Strong 1997). The standard picture on which these multi-component models are based is that cosmic rays are accelerated in sources (probably supernova remnants) in the disk; then they diffuse due to scattering off self-generated magnetohydrodynamic waves over a large region, referred to as the "diffusive cosmic-ray halo," until they reach a boundary at which free escape into the intergalactic medium occurs. Such models can account for locally measured cosmic-ray isotopic ratios (over a wide mass range of both stable and unstable nuclei), as well as diffuse gamma-ray (Strong & Mattox 1996) and radio synchrotron (*e.g.*, Beuermann, Kanbach, & Berkhujsen 1985) emission in our Galaxy (Webber, Lee & Gupta 1992; Strong & Moskalenko 1998; Strong, Moskalenko & Reimer 2000). But halos of arbitrary vertical thickness are not allowed. The best-fit models presented by Webber, Lee & Gupta (1992), for example, "conservatively limit the cosmic-ray halo to a maximum thickness of 4 kpc." This halo size also is consistent with measurements of the extent of radio-synchrotron emission in other, edge-on galaxies (Hummel, Smith & Van de Hulst 1984; Harnett & Reynolds 1985).

Drawing on the results of their detailed numerical diffusion model, Strong & Moskalenko (1998) have previously illustrated how the the cosmic ray spectrum must vary throughout the volume of our Galaxy in order to satisfy this collection of observational constraints. Using expression (2), it is straightforward to integrate over the spectrum of cosmic rays that results from such a model and determine the ionization rate due to cosmic rays at each location in the disk of the Galaxy and throughout the vertically extended diffusive cosmic-ray halo. The result of such an integration is shown here in Fig. 1. The model from which Fig. 1 has been derived does not provide a unique solution for the structure of the cosmic-ray halo, but its features are typical of the best-fit model results that have emerged from different research groups in recent years. The ionization rate (and its associated flux of cosmic rays) shows modest radial variation within the disk of the Galaxy reflecting, for example, a nonuniform source distribution. But the amplitude of this variation decreases with height above the disk plane and, most importantly in the context of this paper, the diffusive halo has a boundary that is a few kpc above and below the disk.

Outside of the volume defined by the surface of the diffusive cosmic-ray halo, the flux of cosmic rays drops rapidly as the high-energy particles stream freely into the general intergalactic medium. We infer from this as well that, outside of the cosmic-ray halo, the ionization rate due to cosmic rays drops significantly below its value in the local interstellar medium. A reasonable upper limit to the mean intergalactic cosmic ray flux and associated intergalactic ionization rate, $\zeta_{\rm IG}$, can be obtained by asking what the dilution factor would be if the cosmic rays, which are generally trapped within the diffusive halo in a volume,

$$V_{\rm halo} \sim \pi R_{\rm disk}^2 Z_{\rm halo} \sim \pi(15\,{\rm kpc})^2(4\,{\rm kpc}) = 3\times10^3\,{\rm kpc}^3\,, \tag{6}$$

for a time $t_{\rm halo} \approx 10^8$yr, were to be spread over the average volume that is occupied by each field galaxy, namely,

$$V_{\rm galaxy} \sim 1\,{\rm Mpc}^3 = 10^9\,{\rm kpc}^3. \tag{7}$$

and to accumulate for a Hubble time, $t_{\rm Hubble} \approx 10^{10}$yr. From the ratio of these volumes and timescales we deduce that,

$$\zeta_{\rm IG} \approx \zeta_{\rm CR}(\frac{V_{\rm halo}}{V_{\rm galaxy}})(\frac{t_{\rm Hubble}}{t_{\rm halo}}) \approx 3\times10^{-4}\zeta_{\rm CR} = 3\times10^{-21}\,{\rm s}^{-1}\,. \tag{8}$$

(A similar dilution factor can be derived by demanding that the flux of particles just inside the surface of the diffusive cosmic-ray halo equal the flux of particles just outside that surface and realizing that the “streaming” velocity inside the halo is typically 10^{-4} c; see §20.4 of Longair 1992).

We suspect therefore that if, through a tidal encounter with another galaxy, a gravitationally bound molecular cloud in a spiral galaxy were to be “lifted” intact to a height $\gtrsim 4$ kpc above the disk, it would find itself in an environment where the ionizing flux due to cosmic rays is several orders of magnitude below its value in the disk. Because, according to eq. (5), this would allow the ambipolar diffusion time to drop below the local free-fall time, this environment would also be extremely conducive to formation of stars on a free-fall timescale. Alternatively, as the general axisymmetric and planar structure of a spiral galaxy disk becomes destroyed during a galaxy-galaxy encounter, it seems quite likely that the large-scale, ordered magnetic field associated with that disk will be stretched out and distorted as well. Because the cosmic rays are strongly coupled to the field, such distortions would likely lead to a drop in the cosmic-ray flux by pure dilution as well as through adiabatic energy losses. This would also promote more efficient star formation in the molecular clouds that remain near the midplane of the original disk. It is these general ideas that lead us to suggest that there is a direct link between the magnetic field structure, cosmic ray flux, and the formation of bound star clusters during galaxy-galaxy interactions.

5. Summary and Conclusions

Giant molecular clouds are the primary sites of star formation in the disk of our Galaxy and, as summarized by Jog & Solomon (1992) and Schweizer & Seitzer (1998), it is tempting to suggest that they are the objects from which massive, young clusters of stars form during strong galaxy-galaxy interactions. In addition to the arguments given by these authors, this idea is also appealing from the standpoint of simplicity. That is, why should galaxies go to the trouble of assimilating new molecular cloud structures from diffuse gas during a collision when such clouds already exist in the pre-collision galaxy disks in sufficient quantity to explain the amount of molecular gas seen in interacting systems (Gao & Solomon 1999)?

It is also fair to say that the GMCs in our Galaxy are conveniently poised to undergo an explosive episode of star formation. Within a typical GMC, the thermal Jeans mass is on the order of a stellar mass which, in turn, is much much smaller than the total cloud mass; and nonlinear amplitude density fluctuations are already present on these small scales. However, within the disk environment of our Galaxy, the GMC "bombs" are prevented from exploding because even their clumpy, overdense regions are partially ionized and therefore coupled to the Galaxy's magnetic field. The energy density in the magnetic field exceeds the thermal energy density (so the thermal Jeans mass is not a particularly relevant scale to discuss in the context of cloud fragmentation) and is sufficient to support the clouds against dynamical collapse. Instead, star formation proceeds on a timescale that is governed by ambipolar diffusion which, given the estimated levels of molecular gas ionization within GMCs, is roughly an order of magnitude longer than the cloud's free-fall time. Hence, star formation proceeds in a relatively inefficient manner and, in the disk of our Galaxy, the GMC "bombs" generally fizzle rather than explode.

Modern models of cosmic-ray propagation clearly indicate, however, that in regions outside the disk of our Galaxy the flux of cosmic rays is significantly lower than it is in the disk. Hence, a gravitationally bound GMC that is displaced from the disk should experience a drop in its fractional ionization and be less tightly coupled to the Galaxy's magnetic field. Outside of the diffusive cosmic-ray halo in particular, we have shown that the ionization rate due to cosmic rays is likely to be three to four orders of magnitude below its value at the midplane of the disk. In such an environment, the ambipolar diffusion time that usually governs the rate of star formation in a GMC should drop below the free-fall time and, as a result, dense clumps throughout a GMC should be able to collapse gravitationally on a free-fall timescale. That is to say, a typical GMC "bomb" should ignite in such an environment. Because star formation would proceed throughout the GMC on the same short timescale, very high star formation efficiencies would be achievable and it should be possible for the GMC to form a massive, bound cluster of stars.

We propose that this is essentially what happens during strong galaxy-galaxy interactions. As they move along more or less ballistic trajectories, GMCs originating in one or both galaxies may find themselves displaced from their parent galaxy disk into an environment that has a much lower flux of ionizing cosmic rays. When this happens, each displaced GMC should explode in a burst of star formation and generate a luminosity per unit mass that is consistent with observed starburst phenomena. Alternatively, as the disk of a colliding galaxy becomes severely distorted through tidal interactions, the galaxy's ordered magnetic field may become sufficiently distorted that the "magnetic bottle" that previously had been responsible for confining cosmic rays to a volume near the disk "opens up" and permits the cosmic rays to escape more freely into the intergalactic environment. Then even GMCs that have remained fairly close to the disk's midplane may experience an environmental change that is sufficient to ignite the GMC "bomb."

Although a reduction in the flux of cosmic rays will promote efficient star formation primarily because it leads directly to a reduction in the ambipolar diffusion timescale in the dense cores of molecular clouds, it may also promote cloud fragmentation and gravitational collapse because it will remove an important source of heating from the cloud. It is also reasonable to expect that an environment where the cosmic ray flux is lower will also be an environment where the flux of ionizing photons is lower. Hence, the ambipolar diffusion timescale may also drop significantly in the lower density regions of a GMC where photoionization generally contributes to observed levels of fractional ionization. If the GMC, as a whole, becomes less tightly coupled to the magnetic field as a result of the reduced flux of cosmic rays, the entire cloud should also be expected to contract, leading to a star cluster whose radius is smaller than the original GMC radius.

We prefer this mechanism for instigating the starburst phenomenon to models that rely upon cloud-cloud collisions because (a) the probability of direct collisions between GMCs during galaxy-galaxy interactions is very low (Jog & Solomon 1992) and (b) as we have explained in an accompanying appendix, generally speaking isothermal cloud-cloud collisions do not promote the process of gravitational fragmentation. A model that relies upon environmental over-pressures to trigger the collapse and fragmentation of existing GMCs during galaxy-galaxy interactions, as proposed by Jog & Solomon (1992), has its merits. But it fails to explain how the GMCs are simultaneously able to decouple from the magnetic field of the galaxies in which they reside. Since, in our view, it is coupling to the magnetic field that prevents the GMC "bombs" from exploding while residing in their parent galaxy disk, we are unconvinced that environmental over-pressures can serve as effective triggers of the starburst phenomenon. Because our proposed mechanism for igniting a GMC "bomb" does not rely upon violent gas-dynamical triggers, it may also explain why systems like the LMC are able to form massive, bound star clusters during relatively mild tidal encounters with a galaxy neighbor.

It is not yet clear to us what observations can be made that will clearly differentiate between our proposed mechanism for igniting the starburst phenomenon and other proposed mechanisms that rely primarily on gas-dynamical triggers. Because cosmic rays originate from supernovae and supernova remnants, and the frequency of supernova events is enhanced in starburst galaxies, one should not be surprised to find that systems presently undergoing a starburst event will have enhanced, rather than suppressed, cosmic ray fluxes. Hence, it seems unlikely that studies of starburst galaxies will provide the necessary discriminating diagnostics. Perhaps it will be possible to deduce the extent of the diffusive cosmic ray halo (through, for example, detection of radio synchrotron emission) as well as the distribution of giant molecular clouds in galaxy pairs that are close, but have not yet actually collided.

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A. Why Isothermal Cloud Collisions are Inefficient Triggers of Star Formation

It is often argued that cloud-cloud collisions are effective triggers of star formation. At first thought such an argument seems reasonable because, during an isothermal collision for example, hydrodynamical processes alone can naturally drive the gas to much higher densities, and at higher densities both the free-fall time and the Jeans mass drop from their values in the initially undisturbed cloud. But a more careful examination of all the timescales and lengthscales that are relevant during isothermal cloud collisions illustrates that the process is an inherently ineffective trigger of gravitational cloud fragmentation and star formation. In what follows we briefly present the line of reasoning that brings us to this conclusion.

For simplicity we consider the head-on collision between two initially uniform, spherical clouds of equal mass density ρ_0 and radius R and, hence, equal mass, $M_{\rm cloud} \approx \pi\rho_0 R^3$. (In the spirit of a back-of-the-envelope calculation, we will not keep track of factors of order unity, such as the factor of 4/3 that should otherwise accompany this determination of $M_{\rm cloud}$.) We assume that both clouds have the same internal temperatures, characterized by a sound speed c, and from this sound speed we can express the mass of each cloud as a fraction of the relevant initial Jeans mass M_J as follows,

$$m_J \equiv \frac{M_{\rm cloud}}{M_J} \approx \pi G\rho_0 R^2/c^2. \tag{A1}$$

We will also assume that $m_J \lesssim 1$ because, if m_J were greater than one, the clouds would initially already be in a state of free-fall collapse so a collision would not be required to trigger star formation. Finally, we assume that the clouds are approaching each other with a center of mass velocity $|v|$, so the collision can be characterized by the Mach number, $\mathcal{M} \equiv v/c$.

The first timescale of interest is the collision time,

$$\tau_{\rm collide} \equiv \frac{R}{v} = \frac{1}{\mathcal{M}} m_J^{1/2} (\pi G\rho_0)^{-1/2}. \tag{A2}$$

This is the time it will take for all of the mass in both clouds to flow through their respective isothermal shock fronts and become incorporated into the high-density, compressed slab that is created by the collision. Over the time interval $0 \le t \le \tau_{\rm collide}$, the mass contained in the slab grows roughly linearly with time as $M_{\rm slab} \sim 2M_{\rm cloud}(t/\tau_{\rm collide})$, so the surface density of the slab material varies with time as,

$$\sigma = \frac{M_{\rm slab}}{\pi R^2} \approx 2\rho_0 R(t/\tau_{\rm collide}), \tag{A3}$$

where we have assumed that the clouds first touch at time $t = 0$.

Now we determine to what mass density $\rho_{\rm slab}$ the clouds become compressed during their head-on collision; at the same time we will determine an approximate half-thickness $Z_{\rm slab}$ of the compressed slab. To do this, we draw upon the Elmegreen & Elmegreen (1978) derivation of the equilibrium properties of a self-gravitating isothermal slab that is under the influence of an external pressure, $P_{\rm ext}$. In our case, the ram pressure of the inflowing gas provides the relevant external pressure, *i.e.*,

$$P_{\rm ext} \approx \rho_0 v^2. \tag{A4}$$

So, from expression (9) of Elmegreen & Elmegreen (1978) we deduce that,

$$\rho_{\rm slab} \approx \frac{1}{c^2}[\rho_0 v^2 + \frac{1}{2}\pi G \sigma^2] \approx \rho_0 \Big[\mathcal{M}^2 + 2 m_J (t/\tau_{\rm collide})^2\Big]. \tag{A5}$$

When discussing cloud-cloud collisions in the context of galaxy interactions it is most realistic to consider the case where the Mach number $\mathcal{M}$ is much greater than one. Then, to a high degree of accuracy, $\rho_{\rm slab}$ is independent of time (over the time interval $0 \leq t \leq \tau_{\rm collide}$) and given by the expression,

$$\rho_{\rm slab} \sim \rho_0 \mathcal{M}^2. \tag{A6}$$

Also, from expressions (4)-(7) of Elmegreen & Elmegreen (1978) we deduce that

$$\frac{Z_{\rm slab}}{R} \approx \frac{1}{\mathcal{M}^2}\,(t/\tau_{\rm collide}). \tag{A7}$$

(This may also be derived by demanding that $M_{\rm slab} = \pi R^2 Z_{\rm slab} \rho_{\rm slab}$.)

With expression (A6) in hand, we can now determine the second timescale of interest, namely, the growth rate — that is, the e-folding time $\tau_{\rm growth}$ — of unstable perturbations in the compressed slab. According to Elmegreen & Elmegreen (1978),

$$\tau_{\rm growth} = \frac{1}{\Omega}(\pi G \rho_{\rm slab})^{-1/2} \approx \frac{1}{\Omega}\frac{1}{\mathcal{M}}(\pi G \rho_0)^{-1/2}, \tag{A8}$$

where the relevant value of the dimensionless frequency Ω depends on the wavelength of the perturbation, but always falls within the range $0 \leq \Omega \lesssim 0.5$. Hence, the e-folding time of perturbations in the slab is never shorter than the collision time $\tau_{collide}$, so if gravitational fragmentation of the slab is to take place, it must happen sometime after the collision is "complete" and the mass of both clouds has been completely incorporated into the slab.

This brings us to the third timescale of interest, that is, the lifetime $\tau_{\rm life}$ of the slab after the collision has been completed. As Elmegreen & Elmegreen (1978) point out, if the compressed layer is destroyed relatively rapidly after its creation, perturbations will not have sufficient time to grow. In our case, a strict upper limit to $\tau_{\rm life}$ is set by the fact that the slab must reexpand into the original cloud volumes after the collision is complete because the ram pressure drops to zero. [We know that it is the ram pressure, rather than the self-gravity of the system, that confines the gas to a slab during the collision because $\mathcal{M}^2$ dominates over the other term in the square brackets of eq. (A5).] Immediately following the collision, in the direction perpendicular to the plane of the slab, a rarefaction wave will propagate from the surface toward the midplane of the slab and the slab will significantly expand on a timescale given by this (vertical) sound crossing time, that is,

$$\tau_{\rm life} \sim \frac{Z_{\rm slab}}{c} \sim \frac{1}{\mathcal{M}^2}\frac{R}{c} \sim \frac{1}{\mathcal{M}^2} m_J^{1/2} (\pi G \rho_0)^{-1/2}. \tag{A9}$$

Hence, the slab lifetime will be shorter by a factor of $1/\mathcal{M}$ than either of the other two relevant timescales.

In summary, because $\tau_{\rm growth}$ is not shorter than $\tau_{\rm collide}$ and $\tau_{\rm growth} \gtrsim \mathcal{M}\tau_{\rm life}$ for all perturbation wavelengths, it is very unlikely that any gravitationally driven fragmentation will occur during an isothermal collision between two gas clouds.

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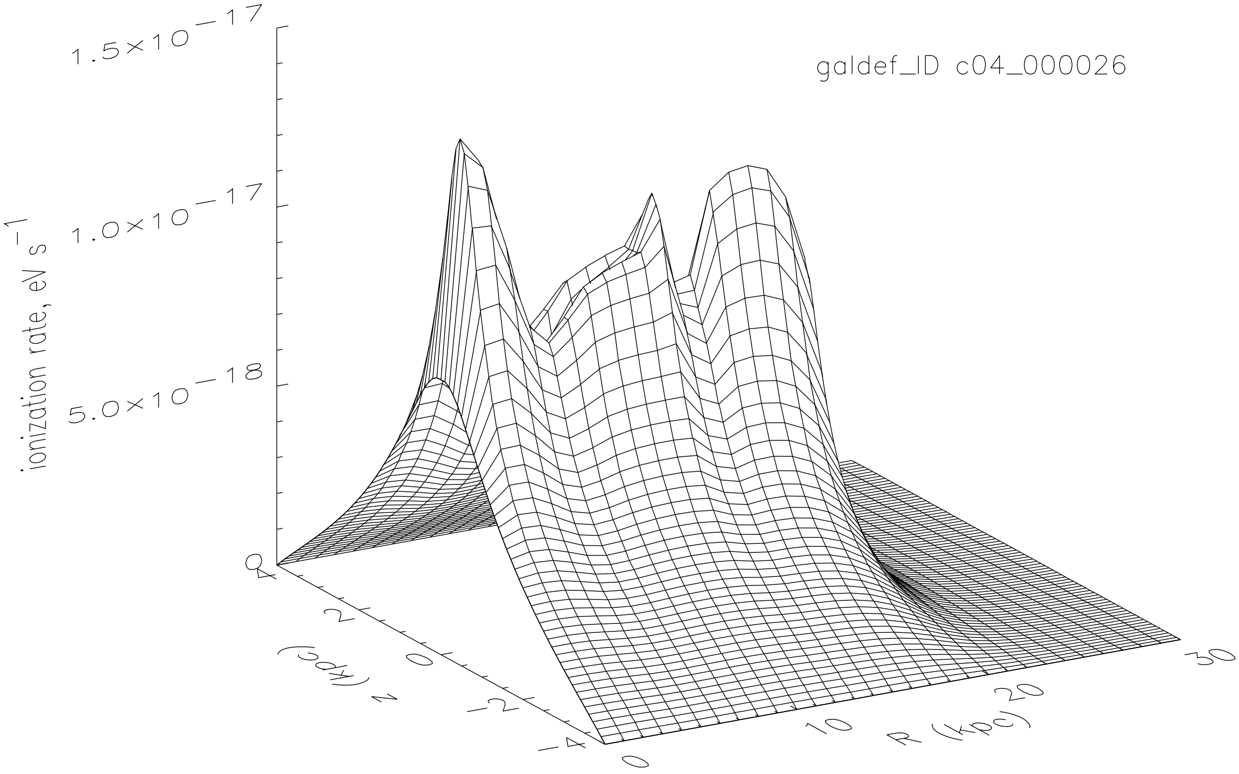
Fig. 1.— The ionization rate due to cosmic rays is shown as a function of radial R and vertical z position in the diffusive cosmic-ray halo that encompasses our Galaxy. The flux of the cosmic-ray protons and Helium from which these results have been derived comes from the detailed numerical diffusion model described by Strong & Moskalenko (1998), and corresponds to case 'HEMN' in that paper. In this model, the Galaxy is assumed to be axisymmetric and the surface of the diffusive cosmic-ray halo is located at 4 kpc above and below the midplane of the Galaxy's disk.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1.5×10^{-17}
1.0×10^{-17}
5.0×10^{-18}
0
z (kpc)
4
2
0
-2
-4
R (kpc)
0
10
20
30