THE CONTRIBUTION OF DIFFERENT SUPERNOVA POPULATIONS TO THE GALACTIC GAMMA-RAY BACKGROUND

E. G. Berezhko
Institute of Cosmophysical Research and Aeronomy, Lenin Avenue 31, 677891 Yakutsk, Russia; berezhko@ikfia.ysn.ru

AND

H. J. Völk
Max-Planck-Institut für Kernphysik, P.O. Box 103980, D-69029 Heidelberg, Germany; Heinrich.Voelk@mpi-hd.mpg.de

Received 2004 March 2; accepted 2004 April 15

ABSTRACT

The contribution of source cosmic rays (SCRs), accelerated and still confined in supernova remnants (SNRs), to the diffuse high-energy $\gamma$-ray emission above 1 GeV from the Galactic disk is studied. The $\gamma$-rays produced by SCRs have a much harder spectrum than those generated by Galactic cosmic rays (GCRs). Extending a previous paper, a simple SNR population synthesis is considered and the inverse Compton emission from the SCR electrons is evaluated in greater detail. The combined spectrum of $\gamma$-ray emission from the Galactic SNR population is then calculated, and this emission at low Galactic latitudes is compared with the diffuse $\gamma$-ray emission observed by EGRET and ground-based instruments. The average contribution of SCRs is comparable to the GCR contribution already at GeV energies, resulting from supernovae of Types II and Ib exploding into the wind bubbles of quite massive progenitor stars, and becomes dominant at $\gamma$-ray energies above 100 GeV. At TeV energies, the dominant contribution is from SCRs in SNRs that expand into a uniform interstellar medium. In fact, the sum of hadronic and inverse Compton $\gamma$-rays would exceed the limits given by the existing experimental data unless the confinement time $T_{SN}$, i.e., the time until which SNRs confine the main fraction of accelerated SCRs, is as small as $T_{SN} \approx 10^4$ yr and the typical magnetic field strength in SNRs as large as 30 $\mu$G. However, both situations are possible as a result of field amplification through cosmic ray (CR) back-reaction in the acceleration process. We point out that accurate measurements of the low-latitude diffuse Galactic $\gamma$-ray spectrum at TeV energies can serve as a unique consistency test for CR origin from the SNR population as a whole.

Subject headings: cosmic rays — diffuse radiation — gamma rays: theory — supernova remnants

1. INTRODUCTION

The diffuse Galactic $\gamma$-ray emission was measured with the Energetic Gamma Ray Experiment Telescope (EGRET) on the Compton Gamma Ray Observatory (CGRO) up to $\approx$20 GeV. To first approximation it can be described by a suitable model for the diffuse interstellar gas, the Galactic cosmic ray (GCR) distribution, and the diffuse photon fields (Hunter et al. 1997b, 1997a). The resulting GCR distribution in the Milky Way essentially corresponds to that measured in situ in the local neighborhood of the solar system. This is also our starting point.

However, above $\approx$1 GeV the observed average diffuse $\gamma$-ray intensity in the inner Galaxy, 300° < $l$ < 60°, |$b$| < 10°, exceeds the model prediction significantly, and this excess appears to increase monotonically with $\gamma$-ray energy. Spatially, the diffuse $\gamma$-ray intensity strongly increases toward the Galactic midplane.

In a previous paper (Berezhko & Völk 2000a, hereafter Paper I), we have calculated the contribution of an unresolved distribution of cosmic ray (CR) sources to the diffuse $\gamma$-ray flux in the disk, assuming the GCR distribution within its large confinement volume to be the same as that measured in situ. The CR sources contribute through the $\gamma$-ray emission of the accelerated particles that are confined in their interiors: the source cosmic rays (SCRs). It was assumed that supernova remnants (SNRs) are the dominant source of GCRs; on this premise, it was found that the SCRs make an important contribution to the diffuse high-energy $\gamma$-ray emission at low Galactic latitudes. Since the CR energy spectrum inside SNRs is much harder than on average in the Galaxy (the average spectrum being softened by rigidity-dependent escape from the Galaxy in the diffusion/convection region above the disk), the relative SNR contribution increases with energy and was in fact found to become dominant at $\gamma$-ray energies $\epsilon_{\gamma} \gtrsim$ 100 GeV. This led to a substantial increase of the “diffuse” $\gamma$-ray emission from the Galactic disk so as to constitute a significant observational background. This hard-spectrum background must also be taken into account in the search for spatially extended Galactic CR sources above the GeV region.

Extending Paper I, we take into account here that a significant fraction of the SNR population, the Type II and Ib supernovae (SNe) with progenitor masses in excess of 15 $M_{\odot}$, expand into a nonuniform circumstellar medium strongly modified by the wind from the progenitor star. In such SNRs, the majority of CRs should be produced in a thin dense shell of swept-up interstellar medium (ISM). As a consequence of the high density of this shell, the expected $\pi^0$-decay $\gamma$-ray emission with energies $\epsilon_{\gamma} \lesssim$ 100 GeV is shown to be an order of magnitude larger than for an identical explosion into a uniform ISM. This provides a natural explanation for the observed diffuse $\gamma$-ray excess for energies $\gtrsim$1 GeV. We also consider in detail the contribution of the high-energy SCR electrons to the high-energy $\gamma$-ray background due to inverse Compton (IC) scattering on the background photon field. Since the energy spectrum of the SCR electrons is strongly influenced by synchrotron losses, the IC $\gamma$-ray emission depends on the effective magnetic field strength in the SNR interior. This field
may in turn be amplified relative to the external field by the acceleration process itself (Bell & Lucek 2001). We take such a field amplification into account as an average property of SNRs relative to the contribution from the GCRs. As demonstrated below, this results in an order of magnitude increase of the Galactic \(\gamma\)-ray emission at TeV energies. In a preliminary form, these results have been presented in Berezhko \\
and Völk (2003).

We therefore conclude that a detailed measurement of the low-latitude diffuse Galactic \(\gamma\)-ray spectrum within the energy interval from 1 to 10^4 GeV will constitute an important indirect test for the GCR origin from the population of Galactic SNRs as a whole.

2. GAMMA-RAY LUMINOSITY OF SNRS EXPANDING INTO A UNIFORM CIRCUMSTELLAR MEDIUM

The majority of the GCRs with charge number \(Z\), at least up to kinetic energies \(\epsilon \sim 10^{15}Z\) eV, are presumably accelerated in SNRs. Individual examples for which this is the case are the objects SN 1006 (Berezhko et al. 2002, 2003a) and Cas A (Berezhko et al. 2003b; Berezhko \\
and Völk 2004). According to theory, a significant part of the hydrodynamic SN explosion energy \(E_{SN} \approx 10^{51}\) ergs is converted into CRs already in the early Sedov phase of the explosion, \(t \sim 10^3\) yr, as a result of diffusive shock acceleration (Berezhko et al. 1996; Berezhko \\
and Völk 1997, 2000b). Later on, the total CR energy content \(E_c\), and therefore the \(\gamma\)-ray production slowly varies with time. We neglect this variation in the energetics in our considerations below, assuming a time-independent SCR energy content \(E_c = 0.1E_{SN}\), consistent with the GCR energy budget requirement.

The total number of active SNRs, \(N_{SN} = \nu_{SN}T_{SN} \approx 300-3000\), is an increasing function of their assumed lifetime \(T_{SN} = 10^4-10^5\) yr, i.e., the time until which they can confine the main fraction of the accelerated particles. Here \(\nu_{SN} \approx 1/30\) yr is the Galactic SN rate. Consequently, we conclude that the population of the oldest SNRs should dominate the total \(\gamma\)-ray luminosity of the ensemble of Galactic SNRs. We therefore consider only SNRs in the Sedov phase of their evolution in our estimate of the total \(\gamma\)-ray luminosity.

Thus, \(\gamma\)-ray producing CRs in the Galaxy are the sum of two basically different populations. The first consists of the ordinary GCRs, and presumably occupies a large Galactic residence volume quasi-uniformly. This residence volume exceeds by far the volume of the gas disk that harbors the CR sources (Ptuskin et al. 1997; Berezinsky \\
and Ginzburg 1990). The second CR population is the SCRs in the localized SNRs.

During the initial, active period of SNR evolution at times \(t \leq T_{SN}\) when the SN shock is relatively strong, the volume occupied by the accelerated CRs practically coincides with the shock volume. During later evolutionary stages, the shock becomes weak and CRs begin to leave the SNR acceleration region. After some period of time, the escaping SCRs become very well mixed with the ambient GCRs. Thus, the controlling factor is the shock strength.

The \(\gamma\)-ray production by the GCRs is quite well studied (Hunter et al. 1997a; Mori 1997). Therefore, we shall concentrate on the relative contribution of the SCR population.

It was shown in Paper I that the total \(\gamma\)-ray spectrum measured from an arbitrary Galactic disk volume is expected to be

\[
\frac{dF_\gamma}{d\epsilon_\gamma} = \frac{dF^{GCR}_\gamma}{d\epsilon_\gamma} \left[1.4 + R(\epsilon_\gamma)\right],
\]

where \(dF^{GCR}_\gamma/d\epsilon_\gamma\) is the \(\pi^0\)-decay \(\gamma\)-ray spectrum due to GCRs and the additional factor 0.4 is introduced to approximately take into account the contribution of the GCR electron component to the diffuse \(\gamma\)-ray emission at GeV energies (Hunter et al. 1997a).

For the ratio \(R(\epsilon_\gamma) = Q^{SCR}/Q^{GCR}\) of the \(\gamma\)-ray production rates due to SCRs and GCRs, we have

\[
R(\epsilon_\gamma) = 0.07\epsilon_\gamma \left(\frac{T_p}{10^3\text{yr}}\right) \left(\frac{\epsilon_\gamma}{1\text{ GeV}}\right)^{0.6} (1 + R_{ep}),
\]

with \(\zeta = N^{SCR}/N^{GCR}\), where \(N^{GCR} = 1\) cm\(^{-3}\) and \(N^{SCR}\) are the gas number density in the Galactic disk and inside SNRs, respectively, \(R_{ep} = Q^{IC}/Q^{pp}\) is the ratio of the IC to \(\pi^0\)-decay \(\gamma\)-ray production rates due to the SCRs. \(T_p\) is the proton SCR confinement time, and the difference between the spectral indices of the GCRs and the SCRs is assumed to be 0.6.

In the case of SNRs expanding into a uniform ISM \(T_p\), which is equal to the confinement time \(T_c\) of the SCR electron component if synchrotron losses are not more restrictive for electrons (see below), is given by the expression

\[
T_p = \min\left[T_{SN}, 10^3(\epsilon/\epsilon_{max})^{-5}\text{yr}\right],
\]

where \(\epsilon_{max} = 10^5Z\) GeV is the maximal energy of CRs that can be accelerated in SNRs, and \(\epsilon_s = 0.1\epsilon_c\), which is roughly valid for the hadronic \(\gamma\)-ray production process. The decrease of the proton confinement time \(T_p\) with energy \(\epsilon\) is due to the diminishing ability of the SNR shock to produce high-energy CRs during the Sedov phase. In fact, the highest CR energy \(\epsilon\) and therefore the highest energy \(\gamma\)-ray production slowly decreases with time as \(\epsilon \propto t^{-1/3}\), which implies the escape of the highest energy CRs from the SNR (Berezhko et al. 1996; Berezhko \\
and Völk 2000a).

Since it was shown in Paper I that electron SCRs give a \(\gamma\)-ray contribution comparable with that of the nuclear SCRs, we consider the electron emission here in greater detail. In the relevant energy range the \(\gamma\)-ray luminosity due to the electron IC scattering on the background photons can be described in the Thompson limit (Longair 1981; Berezinsky \\
and Ginzburg 1990). Since the electron spectrum in a SNR depends significantly on the SNR age \(t\) because of synchrotron losses, we first determine the \(\gamma\)-ray production rate \(dQ^{IC}_\gamma(\epsilon_\gamma)/dN^{e}_{SN}\) of one SNR of age \(t\):

\[
\frac{dQ^{IC}_\gamma(\epsilon_\gamma)}{dN^{e}_{SN}} = \sigma_T e N_{ph} \frac{d\epsilon_e}{d\epsilon_\gamma} \frac{dN^{e}_{SCR}(\epsilon_\gamma)}{dN^{e}_{SN}},
\]

where

\[
\epsilon_e = m_e c^2 \sqrt{3\epsilon_\gamma/(4\epsilon_{ph})}
\]

is the energy of electrons that produce an IC photon with mean energy \(\epsilon_\gamma\). \(\sigma_T = 6.65 \times 10^{-25}\) cm\(^2\) is the Thomson cross section, and \(\epsilon_{ph}\) and \(N_{ph}\) are the mean energy and number density of the background photons, respectively. Finally,

\[
dN^{e}_{SCR}/dN^{e}_{SN} = N_{SCR}^{e}/V_g
\]

denotes the spatial electron SCR number density averaged over the Galactic disk volume \(V_g\), where \(dN^{e}_{SN} = \nu_{SN} dt\) is the number of SNRs of age between \(t\) and \(t + dt\) that contribute to the IC emission at energy \(\epsilon_\gamma\) from CR electron sources.
The source electron spectrum at energies \( \epsilon < \epsilon_t \), where synchrotron losses are not important, can be represented in the form
\[
N_{\text{SCR}}^e(t, \epsilon) = K_{\epsilon p} N_{\text{SCR}}(\epsilon),
\]
where \( K_{\epsilon p} \approx 10^{-2} \) is the electron to proton ratio and \( N_{\text{SCR}}(\epsilon) d\epsilon \) is the total number of SCR protons per SNR in the energy interval \( d\epsilon \). Finally,
\[
\epsilon_t = 1.25 \left( \frac{10^5 \text{ yr}}{t} \right) \left( \frac{10 \mu \text{G}}{B} \right)^2 \text{TeV}
\]
is the low-energy limit of the part of the electron spectrum that is modified by the synchrotron losses, and \( B \) is the magnetic field strength inside the SNRs. For \( \epsilon > \epsilon_t \), the electron spectrum is steeper because of synchrotron losses:
\[
N_{\text{SCR}}^e(t, \epsilon) = K_{\epsilon p} N_{\text{SCR}}(\epsilon)/(\epsilon_t/\epsilon).
\]

We also neglect the time dependence of the proton spectrum in SNRs, since in the later Sedov phase, which is most important here, it is not significant, especially at high energies (Berezhko et al. 1996). The proton spectrum is taken in a power-law form \( N_{\text{SCR}} \propto \epsilon^{-\gamma} \), with \( \gamma = 2.15 \), consistent with the requirements for the sources of the Galactic CRs.

The maximum energy of SCR electrons is also restricted by their synchrotron losses during their acceleration at the SNR shock:
\[
\epsilon_{\text{max}}^e = 24 \left( \frac{V_s}{10^3 \text{ km s}^{-1}} \right) \left( \frac{B}{10 \mu \text{G}} \right)^{-1/2} \text{TeV},
\]
if \( \epsilon_{\text{max}}^e < \epsilon_{\text{max}}^s \); here \( V_s \) is the shock speed. At the beginning of the Sedov phase (\( t \sim 10^3 \) yr), the shock speed \( V_s \approx 4 \times 10^5 \) km s\(^{-1}\) and subsequently decreases as \( V_s \propto t^{-3/5} \). Therefore, the electron confinement time can be written in the form
\[
T_e = \min \left[ \frac{t_p}{\epsilon}, 10^3 (\epsilon / \epsilon_{\text{max}}^s)^{-5/3} \text{ yr} \right].
\]
The total IC \( \gamma \)-ray production rate due to all existing SNRs can be determined by the expression
\[
Q_{\text{IC}}^\gamma = \int_0^{T_e} \frac{dQ_{\text{IC}}^\gamma}{dN_{\text{SN}}^e} dt.
\]
First we normalize \( Q_{\text{IC}}^\gamma \) to the production rate of \( \pi^0 \)-decay \( \gamma \)-rays from inelastic CR-gas collisions, primarily \( p-p \) collisions, which can be written in the form (Drury et al. 1994)
\[
Q_{\text{IC}}(\gamma) = Z_s \sigma_{pp} N_{\gamma} n(\epsilon),
\]
where \( N_{\gamma} \) is the local gas number density, \( \sigma_{pp} \) is the inelastic \( p-p \) cross section, \( Z_s \) is the so-called spectrum-weighted moment of the inelastic cross section, \( n(\epsilon) = N_{\text{SCR}}(\epsilon)/T_p V_0 / T_{\gamma} \) is the spatial number density of SCRs averaged over the Galactic volume \( V_0 \), and \( c \) is the speed of light. Then, performing the integration over all possible SNR ages, we find the ratio of luminosities:
\[
R_{\epsilon p} = 73.6 K_{\epsilon p} \left( \frac{N_{\text{ph}}}{N_{\text{SCR}}} \right) \left( \frac{4 \epsilon / \epsilon_{\text{ph}}}{3m_e^2 c^4} \right)^{\gamma - 1/2},
\]
for \( \gamma < \gamma^* \), and
\[
R_{\epsilon p} = 73.6 K_{\epsilon p} \left( \frac{N_{\text{ph}} T_p}{N_{\text{SCR}} T_{\gamma}} \right) \left( \frac{4 \epsilon / \epsilon_{\text{ph}}}{3m_e^2 c^4} \right)^{\gamma - 1/2}
\times \left( \frac{\epsilon / \epsilon_t}{\epsilon^*_t} \right)^{-1/2} \left( 1 + \ln \frac{\epsilon_t}{\epsilon_t} \right)
\]
for \( \gamma > \gamma^* \), where
\[
\epsilon^*_t \approx 8 \left( \frac{\epsilon_{\text{ph}}}{1 \text{ eV}} \right) \left( \frac{10^5 \text{ yr}}{T_p} \right)^2 \left( \frac{B}{10 \mu \text{G}} \right)^{-4} \text{TeV}
\]
is the energy of \( \gamma \)-rays that are emitted by electrons with a synchrotron loss time equal to \( T_e \). The value \( Z_t = 0.113 \) was used, which corresponds to \( \gamma = 2.15 \) (Drury et al. 1994).

Because of the rather hard SCR electron spectrum, the dominant contribution to the IC radiation comes from collisions with cosmic microwave background (CMB) photons, which are characterized by \( N_{\text{ph}} = 400 \text{ cm}^{-3} \) and \( \epsilon_{\text{ph}} = 6.7 \times 10^{-4} \) eV. The far-infrared radiation (FIR) field with \( N_{\text{ph}} \approx 20 \text{ cm}^{-3} \) and \( \epsilon_{\text{ph}} \approx 0.01 \) eV contributes ~20% for \( \epsilon_t \geq 100 \text{ GeV} \). In the case of the FIR background, the Klein-Nishina cross section must be used instead of the Thomson limit at TeV \( \gamma \)-ray energies. We take account of the Klein-Nishina decrease of the cross section at high energies through multiplying \( \sigma_t \) by the correction factor:
\[
f_{\text{KN}} = \exp \left( -3.5 \sqrt{\frac{\epsilon_t \epsilon_{\text{ph}}}{m_e^2 c^4}} \right),
\]
which approximately describes the reduction effect (Blumenthal & Gould 1970). For \( \epsilon_t = 10 \text{ TeV} \), \( f_{\text{KN}} \approx 0.1 \) for the FIR contribution.

Bremsstrahlung \( \gamma \)-rays play no role in the average \( \gamma \)-ray background above GeV energies if \( K_{\epsilon p} \ll 0.1 \) (Paper I).

In Figure 1 we present a calculated \( \gamma \)-ray spectrum based on the above equations (1)–(17) with \( T_{\gamma} = 10^5 \) yr, \( B = 10 \mu \text{G} \), \( \epsilon_{\text{max}} = 10^5 \) GeV, \( K_{\epsilon p} = 10^{-2} \), and \( N_{\text{SCR}}^\text{IC} = N_{\text{SCR}}^\text{IC} \). A SCR power-law index \( \gamma = 2.15 \) and a 10% efficiency of SCR production in SNR, required to account for the GCRs, were used.

One can see that the SCR contribution becomes dominant for \( \epsilon_t \geq 100 \text{ GeV} \), where the expected \( \gamma \)-ray spectrum becomes extremely hard. At TeV energies the predicted flux exceeds the lowest HEGRA upper limit (Aharonian et al. 2001) by almost a factor of 2. The SCR electron contribution increases with energy \( \epsilon_t \) and exceeds the proton SCR contribution for energies \( \epsilon_t \geq 300 \text{ GeV} \). At \( \epsilon_t \approx 1 \text{ TeV} \), electron SCRs contribute ~60% of the total \( \gamma \)-ray flux. Therefore, at these high energies the expected \( \gamma \)-ray emissivity of SNRs is only weakly dependent on the most relevant parameter, the maximum confinement time \( T_{\gamma} \), since according to expressions (2) and (15) we have for \( \epsilon_t \sim 1 \text{ TeV} \)
\[
dP_{\gamma}/d\epsilon_{\gamma} \propto 1 + 0.3 \ln (T_{\gamma}/10^5 \text{ yr}),
\]
taking into account that at this energy \( T_{\gamma} = T_e = T_{\gamma} \). Even if \( T_{\gamma} \) is as short as \( 3 \times 10^4 \) yr, the expected SCR contribution at TeV energies is therefore only 30% lower and still exceeds the GCR contribution by more than an order of magnitude (see Fig. 1). We emphasize that in the latter case the SCR contribution is entirely due to the electrons: they radiate ~85% of the TeV \( \gamma \)-rays. This extreme case has to be considered as the
lowest value for the SCR contribution to the Galactic $\gamma$-ray background radiation, since there are no other ways to decrease it. This holds in particular also for the average gas number density $N_{\text{SCR}}$ inside SNRs because the IC $\gamma$-ray production rate does not depend on it. Such a short confinement time $T_{\text{SN}} \sim 10^4$ yr (Ptuskin & Zirakashvili 2003) could be realized in a scenario where a strong amplification of the magnetic field $B$ in SNRs is balanced by nonlinear wave-wave interactions such that the wave magnetic power $P(l)$ in the spatial scale $l$ is dissipated in a (minimal) eddy turnover time $\sim \eta(l)$ (Verma et al. 1996). Here $\eta(l) = v_A [P(l)/B^2]^{1/2}$ is the rms value of the plasma mass velocity at scale $l$ and $v_A$ denotes the Alfvén velocity.

We also point out that the restriction of the electron confinement time $T_e$ due to their synchrotron losses becomes significant for $\gamma$-ray energies $\epsilon_\gamma > 1$ TeV and leads to the steepening of the $\gamma$-ray spectrum (see Fig. 1).

According to Figure 1, for energies $\epsilon_\gamma = 0.4 \text{--} 4$ TeV the expected differential $\gamma$-ray energy spectrum from the central part of the Galaxy has the form

$$dF_\gamma/d\epsilon_\gamma = A(\epsilon_\gamma / 1 \text{ TeV})^{-\alpha},$$

with power-law index $\alpha \approx 2.25$ and amplitude $A = (6 \text{--} 8) \times 10^{-9} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ TeV}^{-1}$, depending on the confinement time $T_{\text{SN}}$.

As one can see from Figure 1, the $\gamma$-ray flux expected for $T_{\text{SN}} = 10^3$ yr at $\epsilon_\gamma = 1$ TeV exceeds the HEGRA upper limit (Aharonian et al. 2001) and the preliminary flux measured by the Milagro detector (Fleysher 2003). We emphasize here that we have multiplied the Milagro flux by a factor of 3, since we present the expected $\gamma$-ray flux from the region $|b| \leq 2^\circ$, whereas the Milagro data correspond to $|b| \leq 5^\circ$. Even for $T_{\text{SN}} = 3 \times 10^4$ yr, the expected flux exceeds the Milagro data.

The only physical parameter that strongly influences the IC $\gamma$-ray production rate is the SNR magnetic field $B$: according to equations (15) and (16), $F_\gamma \propto B^{-2}$. A value of $B$ that is significantly higher than $B = 10 \mu G$ can be attributed to field amplification at the shock front because of the strong wave production by the acceleration of CRs far into the nonlinear regime (Bell & Lucek 2001).

In Figure 2 we present the same calculations as in Figure 1, but with a mean SNR magnetic field value $B = 30 \mu G$. Compared with the previous case, the IC $\gamma$-ray emission is decreased by an order of magnitude. Therefore, the contribution of the hadronic SCRs to the $\gamma$-ray flux becomes much more significant: at TeV energies and for $T_{\text{SN}} = 10^5$ yr the $\pi^0$-decay $\gamma$-rays exceed the IC $\gamma$-rays by a factor of 2.5, whereas for $T_{\text{SN}} = 3 \times 10^4$ yr the IC $\gamma$-rays still exceed the $\pi^0$-decay $\gamma$-rays by a factor of 1.3. One can see that in this case the expected $\gamma$-ray flux is below the HEGRA and Tibet upper limits already for a SCRR confinement time $T_{\text{SN}} = 10^3$ yr. The expected $\gamma$-ray spectrum at energies $\epsilon_\gamma = 0.1 \text{--} 10$ TeV is characterized by $\alpha = 2.4$ and $A = (2.4 \text{--} 4) \times 10^{-9} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ TeV}^{-1}$, which also considerably exceeds the GCR contribution. Note that since the maximum energy of SCRs obeys $\epsilon_{\max} \propto B$, it is larger by a factor of 3 compared with the previous case. Taking into account the uncertainty of the Milagro flux, one can conclude that this measurement does not contradict our prediction for the case of the amplified magnetic field. The lower value of the confinement time $T_{\text{SN}} \sim 10^4$ yr might be preferable in this sense.

3. WIND SNRS OF TYPE II AND Ib

The above consideration corresponds to the assumption that all existing Galactic SNRs expand into a uniform circumstellar medium with a density of $N_{\text{GCR}}^G = 1 \text{ cm}^{-3}$. However, for stellar masses in excess of $\sim 15 M_\odot$ the progenitors of Type II and Ib...
SN can strongly modify their environments through their ionizing radiation and very energetic winds (Weaver et al. 1977; Chevalier & Liang 1989). Disregarding large-scale turbulent mixing processes in the sequel, the stellar wind from such a star creates a low-density bubble beyond the termination shock out to a radius $R_{sh}$, bounded by a dense swept-up shell of interstellar material. The radius of the wind bubble can be several tens of parsecs. The very hot dilute bubble contains a small amount of mass, typically a few solar masses or even less, whereas the thin swept-up shell ultimately contains thousands of solar masses. This shell will have a gradual inner boundary where localized mixing of bubble and shell material proceeds. Since the typical SN ejecta mass $M_{ej} \approx 10 M_{\odot}$ for massive progenitors, only a relatively small fraction of the SN explosion energy $E_{SN}$ is transformed into internal gas and CR energy during the SN shock propagation through the bubble, and the main part of SN energy is deposited in the shell.

There are no physical reasons why the CR production process in the shell matter is substantially different from that in the uniform ISM case, apart from the different distribution of ion injection across the SN shock surface due to the magnetic field geometry (Völk et al. 2003). The major difference that plays a role for the $\gamma$-decay $\gamma$-rays is that CRs accelerated during the SN shock propagation through the shell are confined in a medium that is more than a factor of 10 denser than the average ISM in the Galactic disk, which is of density $N_{GCR}$. This may lead to an appreciable increase in the overall $\gamma$-ray production.

In addition, the SN shock evolves differently than in the case of a uniform circumstellar medium. The entire active period of SNR evolution, during which the SN shock is strong enough and effectively produces CRs, takes place within the shell of thickness $L$, which is much smaller than its radius, $L \ll R_{sh}$. This implies that during the entire active evolution the SNR shock size $R_s$ remains nearly constant. The time dependence of the shock speed $V_s$, which is very important for the SCR confinement in SNRs, can then be estimated as follows.

The majority of the progenitor stars with an intense wind are main-sequence stars with initial masses $M_1 > 15 M_{\odot}$ (Abbott 1982). In the mean, during their evolution in the surrounding uniform ISM of gas number density $\rho_0 = m_p N_{GCR}^Q$, they create a bubble of size (Weaver et al. 1977; Chevalier & Liang 1989)

$$ R_{sh} = 0.76(0.5 M V_w^2 / \rho_0)^{1/5}, $$

where $\dot{M}$ is the mass-loss rate of the presupernova star, $V_w$ is the wind speed, and $t_w$ is the duration of the wind period.

In order to determine the SN shock dynamics inside the shell, we model the gas number density distribution in the bubble and in the shell in the form

$$ N_g = N_b + (r/R_{sh})^{(\sigma-1)} N_{sh}, $$

where $N_{sh} = \sigma N_{GCR}$ is the peak number density in the shell, $N_b$ is the gas number density inside the bubble, typically very small compared with the shell density, and $\sigma = N_{sh}/N_0$ is the shell compression ratio.

The mass of the bubble

$$ M_b = (4 \pi R_{sh}^3 / 3) m_p N_b $$

is $3-5 M_{\odot}$, whereas the shell mass

$$ M_{sh} = 4 \pi N_{sh} m_p \int_0^{R_{sh}} dr r^2 (r/R_{sh})^{(\sigma-1)} = (4 \pi R_{sh}^3 / 3) N_{GCR} m_p $$

is a few thousandths of a solar mass. Therefore, during SNR shock propagation through the bubble only a small fraction of its energy is given to gas of stellar origin. The main part of the explosion energy is deposited in the shell.

As in the case of a uniform ISM, the evolution of the SN shock can be represented as a sequence of two stages. The first stage is the free expansion phase, which continues up to the time at which the shock sweeps up a mass equal to the ejecta mass $M_{ej}$. The second phase is analogous to the Sedov phase in the uniform ISM. The only significant difference compared with the uniform ISM case is that at the end of the free expansion phase, when a large part of the SN energy has been given to the swept up gas, the shock volume is much larger.

When the swept-up mass becomes large compared with the sum of bubble mass and $M_{ej}$, the SN shock propagation in the shell medium with density $N_g \propto r^{(\sigma-1)}$ can be shown to approach the following self-similar adiabatic solution:

$$ R_s^2 V_s^2 N_g(R_s) = \text{constant}, $$

which corresponds to an expansion law $V_s = R_0 (t / t_0)^{\nu}$, with $\nu = 2 / (3 \sigma + 2)$. The shock decelerates much more rapidly in this case, $V_s \propto t^{-3/(3 \sigma + 2)}$, than in the case of a uniform ISM, where $V_s \propto t^{-3/5}$. Because of the fact that $3 \sigma > 1$, we have approximately $V_s \propto t^{-1}$. During this period of time the shock size and the amount of the swept-up mass are connected by the relation

$$ R_s = R_{sh}(M / M_{ej})^{1/3 \sigma}. $$

Substituting the SN ejecta mass $M_{ej}$ into this expression, we get the value of the SN shock size that corresponds to the end of free expansion phase,

$$ R_s(t_0) = R_{sh}(M_{sh} / M_{ej})^{-1/3 \sigma}. $$

At this epoch the shock speed is roughly the mean ejecta speed,

$$ V_s \approx V_0 = \sqrt{2 E_{SN} / M_{ej}}. $$

Therefore, the corresponding timescale is $t_0 = R_0 / V_0$.

It follows from the expansion law and from equation (24) that the SNR shock reaches the edge of the shell $R_s(t_f) = R_{sh}$ at time

$$ t_f = t_0 (M_{sh} / M_{ej})^{(3 \sigma + 2) / 6 \sigma}, $$

which for $\sigma \gg 1$ gives

$$ t_f \approx t_0 \sqrt{M_{sh} / M_{ej}}. $$

At about this stage, the SNR shock has lost much of its speed and will come into pressure equilibrium with its environment, i.e., $t_f \approx T_{SN}$.

For the typical values, $\dot{M} = 6 \times 10^{-9} M_{\odot}$ yr$^{-1}$, $V_w = 2500$ km s$^{-1}$, and $t_w = 10^7$ yr, which correspond to a progenitor of mass
$M_I = 15 \, M_\odot$ (Abbott 1982; Chevalier & Liang 1989) and $N_{\text{SCR}}^{\text{max}} = 1 \, \text{cm}^{-3}$, we have $R_{sh} \approx 30 \, \text{pc}$ and $M_{sh} = 3 \times 10^4 \, M_\odot$. For $E_{\text{SN}} = 2 \times 10^{51} \, \text{ergs}$ and $M_{sh} = 5 \, M_\odot$, this gives $t_0 = 5 \times 10^7 \, \text{yr}$ and $t_f \approx 10^7 \, \text{yr}$. The ambient interstellar gas may also be photoionized, which implies an external pressure large enough to stop bubble expansion at a lower radius (Chevalier & Liang 1989). Therefore, we also consider below the case $R_{sh} = 17 \, \text{pc}$, which corresponds to $t_f = 3 \times 10^8 \, \text{yr}$.

According to theory, the expanding SNR shock produces a power-law CR spectrum up to a maximum energy (Berezhko 1996; Berezhko et al. 1996; Berezhko & Völk 1997)

$$\epsilon_m \propto R_s V_s,$$  

which is determined by the radius $R_s$ and speed $V_s$ of the shock. The CRs with the highest energy $\epsilon_{\text{max}}$ are produced at the very beginning of the Sedov phase $t \approx t_0$, when the product $R_s V_s$ has its maximum. Subsequently, the product $R_s V_s$ decreases with time approximately as $t^{-1}$ and the SNR shock produces CRs with progressively lower cutoff energy $\epsilon_{\text{min}}(t) < \epsilon_{\text{max}} = \epsilon_{\text{min}}(t_0)$. During that phase, those CRs that were previously produced with energies $\epsilon_{\text{min}}(t) < \epsilon < \epsilon_{\text{max}}$ have now left the remnant without a significant influence on the SNR shock. Therefore, the CR confinement time has to be taken in the form

$$T_p = \min[1, (t_0/t_f)(\epsilon_{\text{max}}/\epsilon)] \, T_{\text{SN}},$$  

where according to the above consideration $T_{\text{SN}} = t_f$. This relation is analogous to equation (3) in the case of a uniform ISM and leads to an increase of the spectral index of the $\gamma$-ray emission by 1 unit.

In order to estimate the highest energy of accelerated CRs in this case, we use the relation $\epsilon_{\text{max}} \propto R_s V_0 B$. Compared with the case of a uniform ISM, where $R_0 = 4 \, \text{pc}$ for $N_{\text{SCR}}^{\text{max}} = 1 \, \text{cm}^{-3}$, the value of $R_0$ is $\sim 10$ times larger but $V_0$ is the same. The magnetic field $B(t_0)$, on the other hand, is much lower than in the case of a SN Ia. Indeed, as follows from equations (25) and (20), the gas number density at the beginning of the Sedov phase is as small as $N_g(t_0) \approx n_0 M_{sh}/M_{sh} \approx N_{\text{SCR}}^{\text{max}}/300$. The magnetic field in the region of the shell is presumably the adiabatically compressed interstellar field $B_0$. If we approximate the directions of $B_0$ as being isotropic, then its magnitude scales with the gas density as $B \propto N_g^{1/3}$ (see, e.g., Chevalier 1974). This gives a field strength $B(t_0) \approx 0.02 B_0$. Taking into account all the factors considered, we conclude that the maximum CR energy in the case of wind SNe with dense shells is roughly a factor of 10 smaller than in the case of SNe exploding into an uniform ISM. Assuming that the field amplification is proportional to $N_g^{1/2}$ (cf. Bell & Lucek 2001), this amplification does not play a role here, since $N_g(t_0)$ is so low.

In contrast to our previous study (Berezhko & Völk 2000b), in which we considered the CR acceleration by SN shocks expanding into a modified bubble, we consider here the opposite extreme of bubble structure. It is assumed that the magnetic field suppresses the mass and heat transport between the dense shell and the hot bubble (Chevalier & Liang 1989). Therefore, the bubble mass is so small that a significant number of CRs are produced by the SN shock only when it enters the dense shell. Even though it is not clear at the moment which of these two concepts is physically more correct, we consider here the possibility of an unmodified bubble because of its great importance for $\gamma$-ray production.

We note that because of the energy dependence of the confinement time $T_p(\epsilon)$, the SCR contribution to the $\gamma$-ray flux undergoes a spectral break at a $\gamma$-ray energy that corresponds to the SCR energy,

$$\epsilon_{\text{break}} = (t_0/t_f) \epsilon_{\text{max}}.$$  

(32)

Protons with energies $\epsilon < \epsilon_{\text{break}}$ survive up to the final active SNR epoch $t = t_f$ and therefore provide the maximum contribution to the $\gamma$-ray SNR emissivity at $\epsilon_{\text{break}} = 0.1 \epsilon_{\text{break}}$, whereas for $\epsilon \geq \epsilon_{\text{break}}$ the $\gamma$-ray emissivity decreases more steeply and goes to zero at $\epsilon_{\gamma} = 0.1 \epsilon_{\text{max}}$.

The gas number density $N_{\text{SCR}}^{\text{max}}$ in expression (2) is a function of SNR age $t$ for wind SNe, because the gas density $N_g(R_s)$ seen by the SN shock changes during the evolution of the SNR. It then follows from equation (20) that the gas density behind the SN shock, where most of the SGRs are located, can be represented as $N_g(t) = \sigma_{\text{SN}} N_g^0 R_s(t) = \sigma_{\text{SN}} (t/t_f) T_{\text{SN}}^3 N_{\text{sh}}$, where $\sigma_{\text{SN}}$ is the SN shock compression ratio. Therefore, the gas density seen by SGRs of energy $\epsilon$ is

$$N_g^{\text{SCR}} = \sigma_{\text{SN}} (t/t_f) T_{\text{SN}}^3 N_{\text{sh}}.$$  

(33)

Since $T_p \propto \epsilon_0^{-1}$ for $\epsilon_0 > \epsilon_{\text{break}}$, the $\gamma$-ray spectrum produced in the shell is very steep at these high energies, and in equation (2) we have $R(\epsilon_0) \propto \epsilon^{-2.4}$.

We assume that the progenitors of Type II SNe are stars more massive than $M_I > 8 \, M_\odot$, and that only those with $M_I > 15 \, M_\odot$ have strong winds that produce extended bubbles, (see, e.g., Abbott 1982). This is also true for the progenitor population of Type Ib SNe. According to Güsten & Mezger (1983), the initial mass function has the form $dn/dM_I \propto M_I^{-\alpha}$, with $\alpha = 1.6, 2.4, 3.24, 3.62$, for the mass intervals $M_I < 1 \, M_\odot$, $1 < M_I < 10 \, M_\odot$, $10 < M_I < 50 \, M_\odot$, and $50 \, M_\odot < M_I$, respectively. Therefore, the ratio of the number of stars whose initial mass $M_I$ exceeds $15 \, M_\odot$ to those with $M_I > 8 \, M_\odot$ is 0.23. Since $\sim 85\%$ of all SN explosions in the Galaxy are Type II and Type Ib SNe (Tammann et al. 1994), and since in addition the explosion energies $E_{\text{SN}}$ of Type II and Type Ib SNe with $M_I > 15 \, M_\odot$ are $\sim 2$ times larger than on average (Chevalier 1977; Hamuy 2003), we conclude that roughly a fraction $\delta = 0.3$ of the total Galactic SN energy release is from wind SNe with progenitors of masses $M_I > 15 \, M_\odot$, which produce extended bubbles.

Therefore, the expected $\gamma$-ray SNR luminosity should be weighted between the SNR populations expanding into a uniform ISM and into bubbles as follows:

$$R = (1 - \delta) R_I + \delta R_{II},$$  

(34)

where $R_I$ corresponds to the ratio of $\gamma$-ray production rates due to SGRs and GCRs if all SNe were exploding into a uniform ISM, and $R_{II}$ is the corresponding ratio for explosions into bubbles.

Assuming that wind SNe have rarefied bubbles with dense shells, the spectrum of $\gamma$-rays calculated with the above ratio $R$, $\sigma = 10$, $\sigma_{\text{SN}} = 5$, and $\delta = 0.3$ is presented in Figures 1 and 2 by the solid lines, which correspond to two different assumptions about the value of the SCR confinement time $T_{\text{SN}}$. It is seen that the wind SNe dominate at energies $\epsilon_0 < 100 \, \text{GeV}$ and fit the EGRET data for $T_{\text{SN}} = 3 \times 10^4 \, \text{yr}$ fairly well, whereas for the larger confinement times $T_{\text{SN}} = 10^5 \, \text{yr}$ their contribution becomes so large that their calculated flux exceeds the EGRET data by a factor of $\sim 2$.

We emphasize that the contribution from the population of wind SNe is mainly due to the nuclear SCR component. The
break in the γ-ray spectra presented in Figures 1 and 2 is at \( \epsilon_0^{\text{break}} = 30 \text{ GeV} \) for \( T_{\text{SN}} = 10^3 \text{ yr} \) and at \( \epsilon_0^{\text{break}} = 100 \text{ GeV} \) for \( T_{\text{SN}} = 3 \times 10^4 \text{ yr} \). For larger energies the proton confinement time and the gas density become lower, \( T_p \propto \epsilon^{-1} \) and \( N_{\text{GR}} \propto \epsilon^{-2} \), which leads to a corresponding decrease of the wind SN contribution to the γ-ray spectrum. This contribution becomes insignificant for \( \epsilon_0 \sim 1 \text{ TeV} \).

4. DISCUSSION

Our investigation is based on a picture in which the population of SNRs, considered here as the main GCR sources, are distributed across the Galactic disk in a manner similar to the interstellar gas. In this case the calculated ratio \( R(\epsilon_0) \) of the SCR to the GCR contribution does not depend on the observing direction in the disk. Since the number of these sources \( N_{\text{SN}} = \nu_{\text{SN}} T_{\text{SN}} \) is very limited, their mean number within the field of view around a few degrees of a stereoscopic system of imaging atmospheric Čerenkov telescopes for γ-ray energies \( \geq 100 \text{ GeV} \) is so low that one should expect large fluctuations of the actual value of \( R(\epsilon_0) \), especially for lines of sight other than those directed toward the inner Galaxy.

The actual SNR distribution within the Galactic disk most probably is not uniform. According to Case & Bhattacharya (1998), the SN explosion rate as a function of galactocentric radius \( r \) has a peak at \( r \approx 5 \text{ kpc} \) and drops exponentially with a scale length of \( \sim 7 \text{ kpc} \). Within \( 5 < r < 20 \text{ kpc} \), this agrees fairly well with the radial distributions of supernovae in a sample of 36 external galaxies of Hubble type Sb–Sbc, which are statistically considered equivalent to our Galaxy (Dragicevic et al. 1999). From our results, for γ-ray energies larger than \( 10 \text{ GeV} \) the ratio \( R \), corrected for the nonuniform SNR distribution, is given by the ratio \( R = \eta(r) R^0 \), where the parameter \( \eta(r) \) represents the ratio of the actual SNR number along the given line of sight to the SNR number that corresponds to their uniform distribution. For lower γ-ray energies, \( 10^0 < \epsilon_0 < 10^4 \text{ MeV} \), where the truly diffuse emission from the GCRs dominates, it is known that the γ-ray emissivity gradient is significantly shallower than the gradient of \( \eta(r) \), as shown by the EGRET instrument on board CGRO (Strong & Mattox 1996), in basic agreement with the results from the COS B satellite (Strong et al. 1988). This much shallower, indeed truly diffuse galactocentric γ-ray gradient can be understood as a nonlinear propagation effect from the disk into the Galactic wind in which the latter is driven by the GCRs themselves (Breitschwerdt et al. 2002). In contrast, the radial galactocentric variation of the direct radiation from the sources is not smoothed by any propagation effect. Therefore, its amplitude should spatially vary in common with the distribution of the sources, integrated along the line of sight, and this variation should become directly measurable in the \( > 10^2 \text{ GeV} \) range. Given several such line-of-sight integrals, one can infer the radial galactocentric γ-ray emissivity gradient. The excellent angular resolution of atmospheric Čerenkov telescopes and the high sensitivity of instruments such as the CANGAROO III and HESS arrays in the Southern Hemisphere, which at TeV energies have an order of magnitude higher sensitivity compared with previous generation instruments such as HEGRA (Aharonian et al. 2001), should allow this measurement, in particular in the inner Galaxy.

The assumed average value \( B = 30 \mu G \) of the magnetic field inside SNRs is considerably smaller than the value of \( 120 \mu G \) inferred for, e.g., SN 1006 in the very early Sedov phase (Berezhko et al. 2002, 2003a). However, the amplification must be a monotonically increasing function of the shock strength. In particular, Bell & Lucek (2001) argue for a dependence \( B \propto V_s^2 \). If this is so, then the effective amplified magnetic field in SNRs varies from \( B \approx 100 \mu G \) at the beginning of the Sedov phase to the typical interstellar value \( B \approx 10 \mu G \) at the end of its active period \( t = T_{\text{SN}} \). Since the value of the magnetic field in SNRs has a dominant influence on the IC γ-ray spectrum, and since all evolutionary phases of a SNR contribute to this spectrum in a roughly equal manner, the adopted value \( B = 30 \mu G \) appears to be quite realistic, averaging over the entire active period.

One should also note that the SNR magnetic field influences the maximum energy of SCRs. Since the most energetic SCRs are created at the beginning of the Sedov phase, their actual maximum energy \( \epsilon_{\text{max}} \propto B(0) \) is expected to be larger than considered here, because the amplified magnetic field \( B(0) \) in this phase is larger than on average by a factor of several. Compared with Figure 2, this effect will produce a more extended and smooth high-energy tail of the γ-ray spectrum up to the energy \( \epsilon_0 \approx 10^3 \text{ GeV} \).

The actual process of SN shock interaction with the dense shell could be much more complicated compared with the ideal picture considered here. If because of some physical factors (inhomogeneities in the surrounding ISM, instabilities, etc.) the shell is always strongly distorted, then one would expect that the SN shock will penetrate through the shell more rapidly. This will lead to a decrease of the propagation time \( t_f \) and/or to the decrease of the effective gas number density \( N_{\text{GR}} \), which in turn will decrease the amount of γ-rays produced in such types of SNRs. In other words, the calculated γ-ray spectrum at \( \epsilon_0 \lesssim 100 \text{ GeV} \) (see Figs. 1 and 2), which is mainly due to the wind SNe, has to be considered as an upper limit, which corresponds to the case of stable unmodified bubbles.

5. SUMMARY

Our considerations demonstrate that the SCRs inevitably make a strong contribution to the diffuse γ-ray flux from the Galactic disk at all energies above a few GeV if the population of SNRs is the main source of the GCRs. According to our estimates, the SCR contribution dominates at energies greater than \( 100 \text{ GeV} \) because of its substantially harder spectrum.

There are two physical parameters that influence the expected γ-ray emission from SNRs significantly: the CR confinement time and the mean magnetic field strength \( B \) inside the SNRs. For a conventional value \( B = 10 \mu G \), the expected γ-ray flux from SNRs exceeds the HEGRA upper limit considerably if the SCR confinement time is as large as \( T_{\text{SN}} = 10^5 \text{ yr} \). This contradiction can be resolved either if we suggest an appreciably higher postshock magnetic field \( B \geq 30 \mu G \) or if the SCR confinement time is as small as \( T_{\text{SN}} = 10^3 \text{ yr} \), or a combination of both effects. These possibilities can be attributed to field amplification by the SCRs themselves. In fact, nonlinear field amplification may also lead to a substantial decrease of the SCR confinement time: according to Ptuskin & Zirakashvili (2003), maximal turbulent Alfvén wave damping with its corresponding increase of CR mobility could make the SCR confinement time as small as \( T_{\text{SN}} = 10^4 \text{ yr} \). Under these circumstances, the most realistic γ-ray background spectrum is represented by the thick solid line in Figure 2. One can see that even in the case of \( B = 30 \mu G \) and \( T_{\text{SN}} \sim 3 \times 10^4 \text{ yr} \), the SCR contribution at TeV energies, with roughly equal numbers of IC and π0-decay γ-rays, still exceeds the GCR contribution by almost an order of magnitude. Note that the preliminary γ-ray
flux at $\epsilon_{\gamma} = 1$ TeV, measured by the Milagro detector (Fleysher 2003), confirms our earlier prediction (Berezhko & Völk 2003).

At lower energies, $\epsilon_{\gamma} \lesssim 100$ GeV, the $\gamma$-ray emission from SNRs is dominated by the wind SNe with initial progenitor mass $M_{\star} > 15 M_\odot$, which expand into the bubble created by the progenitor’s wind, under the assumption that the bubble is not modified by global heat and mass transport. The main fraction of the CRs is produced in this case in which the SN shock propagates through the thin dense shell at the edge of the bubble. Because of the significantly higher mass density of the shell material compared with the mean gas density in the Ga- lactic disk, $\pi^0$-decay $\gamma$-rays dominate at $\gamma$-ray energies $\epsilon_{\gamma} \lesssim 100$ GeV, despite the fact that only $\sim 20\%$ of all SNRs belong to this class of object. As shown in Figures 1 and 2, the discrepancy between the observed diffuse intensity and standard model predictions at energies above a few GeV can be attributed to the SCR contribution. This requires unmodified bubbles for the wind SNe. Since we cannot prove this assumption, the corresponding explanation of the lower energy $\gamma$-ray excess is not a definitive conclusion but rather a plausible suggestion that must await more detailed studies of wind bubble morphologies.

We conclude that a detailed measurement of the low-latitude diffuse Galactic $\gamma$-ray spectrum within the energy interval from $1$ to $10^4$ GeV will allow a strong consistency check for the predominant origin of the GCRs from the Galactic population of SNRs as a whole. In the energy range above a few $100$ GeV our predictions are quite robust, resulting in a hard spectrum whose Galactocentric variation should correspond to that of the observed SNR distribution. If the preliminary Milagro data, which agree with our prediction, are confirmed by more precise measurements, then such measurements will give an indirect confirmation of SNRs as the main sources of GCRs. The measurements will also provide a new tool to study the spatial distribution of SNRs in the Galaxy.

This work has been supported in part by the Russian Foundation of Basic Research, grant 03-02-16524, and by LSS grant 422.2003.2. We thank G. Pühlhofer for a discussion on the observations of diffuse $\gamma$-rays. E. G. B. acknowledges the hospitality of the Max-Plank-Institut für Kernphysik, where part of this work was carried out.

REFERENCES

Abbott, D. C. 1982, ApJ, 263, 723
Aharonian, F. A., et al. 2001, A&A, 375, 1008
Amenomori, M., et al. 2003, in Proc. 28th Int. Cosmic Ray Conf. (Tsukuba), 2305
Bell, A. R., & Lucek, S. G. 2001, MNRAS, 321, 433
Berezhko, E. G. 1996, Astropart. Phys., 5, 367
Berezhko, E. G., Ksenofonov, L. T., & Völk, H. J. 2002, A&A, 395, 943
———. 2003a, A&A, 412, L11
Berezhko, E. G., Pühlhofer, G., & Völk, H. J. 2003b, A&A, 400, 971
Berezhko, E. G., & Völk, H. J. 1997, Astropart. Phys., 7, 183
———. 2000a, ApJ, 540, 923 (Paper I)
———. 2000b, A&A, 357, 283
———. 2003, in Proc. 28th Int. Cosmic Ray Conf. (Tsukuba), 2433
———. 2004, A&A, 419, L27
Berezhko, E. G., Yelshin, V. K., & Ksenofontov, L. T. 1996, J. Exp. Theor. Phys., 82, 1
Berezinsky, V. S., & Ginzburg, V. L. 1990, Astrophysics of Cosmic Rays (Amsterdam: North-Holland)
Blumenthal, G. R., & Gould, R. J. 1970, Rev. Mod. Phys., 72, 237
Breitschwerdt, D., Dogiel, V. A., & Völk, H. J. 2002, A&A, 385, 216
Case, G. L., & Bhattacharya, D. 1998, ApJ, 504, 761
Chevalier, R. A. 1974, ApJ, 188, 501
———. 1977, ARA&A, 15, 175
Chevalier, R., & Liang, P. 1989, ApJ, 344, 332

We conclude that a detailed measurement of the low-latitude diffuse Galactic $\gamma$-ray spectrum within the energy interval from $1$ to $10^4$ GeV will allow a strong consistency check for the predominant origin of the GCRs from the Galactic population of SNRs as a whole. In the energy range above a few $100$ GeV our predictions are quite robust, resulting in a hard spectrum whose Galactocentric variation should correspond to that of the observed SNR distribution. If the preliminary Milagro data, which agree with our prediction, are confirmed by more precise measurements, then such measurements will give an indirect confirmation of SNRs as the main sources of GCRs. The measurements will also provide a new tool to study the spatial distribution of SNRs in the Galaxy.

This work has been supported in part by the Russian Foundation of Basic Research, grant 03-02-16524, and by LSS grant 422.2003.2. We thank G. Pühlhofer for a discussion on the observations of diffuse $\gamma$-rays. E. G. B. acknowledges the hospitality of the Max-Plank-Institut für Kernphysik, where part of this work was carried out.

REFERENCES

Abbott, D. C. 1982, ApJ, 263, 723
Aharonian, F. A., et al. 2001, A&A, 375, 1008
Amenomori, M., et al. 2003, in Proc. 28th Int. Cosmic Ray Conf. (Tsukuba), 2305
Bell, A. R., & Lucek, S. G. 2001, MNRAS, 321, 433
Berezhko, E. G. 1996, Astropart. Phys., 5, 367
Berezhko, E. G., Ksenofonov, L. T., & Völk, H. J. 2002, A&A, 395, 943
———. 2003a, A&A, 412, L11
Berezhko, E. G., Pühlhofer, G., & Völk, H. J. 2003b, A&A, 400, 971
Berezhko, E. G., & Völk, H. J. 1997, Astropart. Phys., 7, 183
———. 2000a, ApJ, 540, 923 (Paper I)
———. 2000b, A&A, 357, 283
———. 2003, in Proc. 28th Int. Cosmic Ray Conf. (Tsukuba), 2433
———. 2004, A&A, 419, L27
Berezhko, E. G., Yelshin, V. K., & Ksenofontov, L. T. 1996, J. Exp. Theor. Phys., 82, 1
Berezinsky, V. S., & Ginzburg, V. L. 1990, Astrophysics of Cosmic Rays (Amsterdam: North-Holland)
Blumenthal, G. R., & Gould, R. J. 1970, Rev. Mod. Phys., 72, 237
Breitschwerdt, D., Dogiel, V. A., & Völk, H. J. 2002, A&A, 385, 216
Case, G. L., & Bhattacharya, D. 1998, ApJ, 504, 761
Chevalier, R. A. 1974, ApJ, 188, 501
———. 1977, ARA&A, 15, 175
Chevalier, R., & Liang, P. 1989, ApJ, 344, 332