

Magnetic Turbulence Boosts Supernova Signals of Axion-Photon Conversion

Damiano F. G. Fiorillo^{1,2}, Ángel Gil Muyor^{1b,3,4}, Georg G. Raffelt^{1b,5}, and Edoardo Vitagliano^{1b,3,4}
¹*Istituto Nazionale di Fisica Nucleare (INFN), Sezione di Napoli, Complesso Universitario di Monte Sant'Angelo, Via Cintia, 80126 Napoli, Italy*
²*Gran Sasso Science Institute (GSSI), L'Aquila, Italy*
³*Dipartimento di Fisica e Astronomia, Università degli Studi di Padova, Via Marzolo 8, 35131 Padova, Italy*
⁴*Istituto Nazionale di Fisica Nucleare (INFN), Sezione di Padova, Via Marzolo 8, 35131 Padova, Italy*
⁵*Max-Planck-Institut für Physik, Boltzmannstraße 8, 85748 Garching, Germany*



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Magnetic fields between a supernova (SN) and Earth convert axions into gamma rays. The absence of such a signal in coincidence with SN 1987A neutrinos, using the coherent Milky Way field, provides well-studied constraints on $g_{ap} \times g_{a\gamma}$ (axion-proton times axion-photon couplings) and on $g_{a\gamma}$ alone. We show that the small-scale power of the turbulent magnetic field component boosts axion-photon conversion and, crucially, extends sensitivity to larger masses. The turbulent field components of the Milky Way and of the Large Magellanic Cloud (hosting SN 1987A) yield improvements of up to 2 orders of magnitude in $g_{ap} \times g_{a\gamma}$. Turbulence should likely impact the sensitivity of other searches based on other axion-photon conversion sites, such as starburst galaxies.

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Introduction—Axions are among the best motivated new particles because they could explain CP conservation in QCD [1–9] and also could be the dark matter of the universe [10–13]. Other axionlike pseudoscalar bosons are predicted by string theory compactifications, realizing the so-called axiverse [14–20]. The smallness of the axion mass, m_a , opens search strategies with low-energy techniques and astrophysical methods. The generic axion–two-photon interaction, $g_{a\gamma} a \mathbf{E} \cdot \mathbf{B}$, enables $a \rightarrow \gamma$ conversion in macroscopic B fields [21,22], which is the basis for most systematic axion searches [23–27] and, through astrophysical B fields, for many constraints [28], and possible future detection, e.g., by SN γ rays [29–32].

Axions interact extremely feebly, so the B -field-induced conversion probability, $P_{a\gamma} = (g_{a\gamma} B L / 2)^2$, is strong only if the path length L is large. However, $P_{a\gamma}$ drops quickly when $L \gg k_a^{-1}$, where $k_a = m_a^2 / 2E_a$ is the axion-photon momentum mismatch. The resulting “loss of coherence” in $a \rightarrow \gamma$ conversion bedevils all axion searches: the conversion volume or path length should be large to overcome the smallness of $g_{a\gamma}$ but is limited by k_a^{-1} . One way to reach larger masses is to give photons an effective $m_\gamma \simeq m_a$ through a dispersive medium [33]. Alternatively, spatial

B -field variations with strong Fourier power at k_a would have a similar effect [34].

In this Letter, we show that the turbulent component of the Milky Way (MW) B field, in the manner of a magnetic undulator, opens a new m_a range, notably for converting axions from the historical SN 1987A. This peculiar Type II SN resulted from the collapse of the blue supergiant Sanduleak –69202 in the Large Magellanic Cloud (LMC) at a distance of 51.4 kpc [35,36] (see Ref. [37] for a slightly smaller distance). The unique neutrino observations [38–45] provided many particle constraints through the duration of neutrino cooling [46–48], notably $g_{aN} \lesssim 10^{-9}$ for the axion-nucleon coupling [31,49–55]. In addition, the absence of a contemporaneous γ -ray signal in the Solar Maximum Mission (SMM) satellite [56–59] provided constraints on $a \rightarrow \gamma$ conversion in the MW B field [60–63]. Assuming axion production through their proton coupling in $np \rightarrow n p a$ bremsstrahlung, the latest bounds on $g_{ap} \times g_{a\gamma}$ [31] are shown in Fig. 1. Conversion in the presumed progenitor field was recently added [30,31], extending sensitivity to larger masses due to the smaller spatial scale of B -field variation. For self-consistency, we have recomputed these constraints with the axion flux spectrum adopted here (see below).

In a similar vein, we here include another previously ignored component, the turbulent MW field. The impact of B -field stochasticity or turbulence on astrophysical $a \rightarrow \gamma$ conversion was previously studied in the context of TeV gamma rays [64–67]. For SN axions, the relevance of a turbulent MW field was proposed in Ref. [29], mostly focusing on a possible boost (up to a factor of 2) compared

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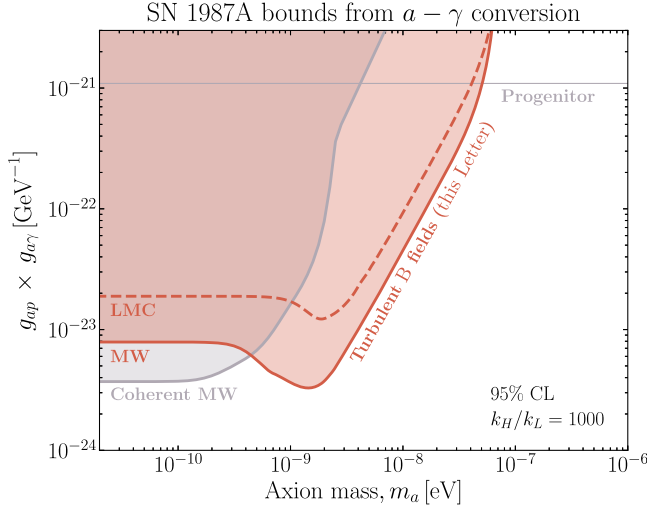


FIG. 1. New constraints (red) on $g_{a\gamma} \times g_{a\gamma}$ from nondetection of SN 1987A gamma rays. We assume $np \rightarrow npa$ bremsstrahlung production and $a \rightarrow \gamma$ conversion in the turbulent B fields of the MW and LMC with $k_H/k_L = 1000$ for the ratio of high and low turbulence scales. In gray, we show previous SN 1987A constraints, labeled “Coherent MW” and “Progenitor,” resulting, respectively, from conversions in the coherent MW field and in the putative dipole field of Sanduleak –69202, the progenitor of SN 1987A (assumed radius $30 R_\odot$ and surface magnetic field 30 G) [31].

to the regular field and the spectral irregularities that would appear in a γ -ray signal from the next Galactic SN for massive axions.

We show that the gains from including turbulent fields in the conversion go far beyond this effect. For the first time, we derive constraints from SN 1987A including the turbulent fields of the MW and of the LMC, separately. For large axion masses, the bounds on $g_{a\gamma} \times g_{a\gamma}$ improve by up to 2 orders of magnitude for neV-range masses, as the turbulent magnetic field provides small-scale structure that facilitates resonant a - γ conversion.

Axion-photon conversion—In all cases of interest here, the conversion probability is small, $P_{a\gamma} \ll 1$, and back-conversion is irrelevant. In this perturbative regime, the conversion amplitude to a photon with polarization j is, up to a global phase, the line of sight (LoS) integral

$$\mathcal{A}_{a\gamma_j} = \frac{g_{a\gamma}}{2} \int_0^L ds B_j(s) e^{-ik_a s}, \quad (1)$$

where $k_a = m_a^2/2E$, the photon plasma mass can be neglected, and $B_j(s)$ are the transverse field components. The conversion probability, summed over both polarizations, then is

$$P_{a\gamma}(k_a) = \frac{g_{a\gamma}^2}{4} \int_0^L ds \int_0^L ds' \mathbf{B}_\perp(s) \cdot \mathbf{B}_\perp(s') \cos[k_a(s-s')], \quad (2)$$

where $\mathbf{B}_\perp$ is the field transverse to the LoS. We denote with $P_{a\gamma}(k_a)$ the conversion probability of an axion with momentum mismatch k_a , and with $P_{a\gamma}$, implicitly depending on m_a , a spectral average. Both quantities coincide for $m_a = 0$. Moreover, we use $\langle P_{a\gamma}(k_a) \rangle$ for an average over B -field configurations. Axions from SN 1987A traveled through one specific realization of the turbulent B field. The conversion probability $P_{a\gamma}(k_a)$ is one draw from a distribution with average $\langle P_{a\gamma}(k_a) \rangle$.

For massless axions, $P_{a\gamma} = (g_{a\gamma}/2)^2 |\int ds \mathbf{B}_\perp(s)|^2$. For the LoS to SN 1987A, the result [31] is given in Table I, derived from the baseline case of the latest coherent MW B -field models of Unger and Farrar [68]. The SN has Galactic coordinates $b_{\text{SN}} = -31.8^\circ$ and $l_{\text{SN}} = 279.7^\circ$, high above the plane. The $B_j(s)$ components along the LoS are shown in Fig. 6 of Ref. [31]. The transverse parts are somewhat weak relative to the longitudinal one, and sizeable fields occur only for about 20 kpc near the Galactic plane. The same $P_{a\gamma}$ descends from a uniform field with $BL = 3.0 \mu\text{G kpc}$.

For m_a so large that $k_a^{-1} \ll L$, there are many a - γ oscillations within L so that $\mathcal{A}_{a\gamma}$ mostly cancels. For $L = 10$ kpc, this occurs for $k_a \gg 6 \times 10^{-28} \text{ eV}$, which for $E = 100 \text{ MeV}$ implies $m_a \gg 0.4 \text{ neV}$. For a uniform field, one finds $P_{a\gamma}(k_a) = (g_{a\gamma} B)^2 k_a^{-2} \sin^2(k_a L/2)$, which for $k_a \rightarrow 0$ returns to $(g_{a\gamma} BL/2)^2$, while for $k_a \gg L^{-1}$, $P_{a\gamma}(k_a)$ oscillates around $(g_{a\gamma} B)^2 k_a^{-2}/2$. For inhomogeneous fields, the same k_a^{-2} scaling is obtained when k_a^{-1} is far below any typical scale of B -field variation.

Turbulent fields—The MW has a strong turbulent field (Table I), mostly inferred from the excess of synchrotron radiation on top of the one expected from the regular field measured through Faraday rotation. It was obtained in Ref. [69] and later updated with *Planck* observations [70], where it was modeled as a Gaussian random field with a Kolmogorov power spectrum and an outer scale equal to $2\pi/k_L = 100 \text{ pc}$. This value is consistent with that inferred

TABLE I. Adopted B -field parameters and $a \rightarrow \gamma$ conversion probability, $P_{a\gamma}$, for massless axions with $g_{a\gamma} = 10^{-11} \text{ GeV}^{-1}$. The turbulent case uses Eq. (3). The coherent MW case uses an integration over the actual profile (see text), whereas the stated $|B|$ and L were chosen to provide the same BL and the same high- k_a suppression of a uniform field.

	$ B , B_0$ [μG]	k_L^{-1} [pc]	L [kpc]	$P_{a\gamma}, \langle P_{a\gamma}(k_a) \rangle$
Milky Way				
Coherent [68]	0.15	...	20	2.1×10^{-3}
Turbulent [69,70]	5	250	4	3.7×10^{-3}
Large Magellanic Cloud				
Coherent [71,72]	1	...	1	2.3×10^{-4}
Turbulent [73]	5	200	0.8	5.9×10^{-4}

in the halo from the fluctuations in the Faraday rotation measure [74,75], although it could be lower in the spiral arms [76]; this is not relevant for this work due to SN 1987A high Galactic latitude. We do not include the striated disk field, introduced in Refs. [69,70], to boost synchrotron radiation compared to the field implied by Faraday rotation. The striated field may not be needed—its effect can also be mimicked by a correlation between cosmic-ray density and B field [69] or the effect of the Local Bubble [77,78]. In any case, for the high-latitude SN 1987A, it is likely unimportant.

For a - γ conversion, a B -field undulating with a given wave number k_B resonantly enhances $P_{a\gamma}$ for $k_a \simeq k_B$. If the Fourier modes are randomly oriented, in projection on the LoS, the resonance condition is met for any $0 \leq k_a \leq k_B$, leading to a nearly flat $\langle P_{a\gamma}(k_a) \rangle$ for $k_a \lesssim k_B$, and to the usual k_a^{-2} variation for $k_a \gg k_B$. The average is taken over field configurations.

The turbulent field is modeled as an isotropic random collection of Fourier modes with a power-law spectrum $|B_k|^2 \propto k^{-\beta-2}$ in a range $k_L \lesssim k \lesssim k_H$ and otherwise zero. The dissipation scale k_H cannot be determined observationally—we show in Supplemental Material [79] that $k_H/k_L = 1000$ is a conservative analytical estimate, to be used in our analysis, while observations hint at much larger values [85]. A Kolmogorov spectrum has $\beta = 5/3$; Goldreich-Sridhar turbulence [86] has an analogous spectrum for the wave number perpendicular to the local magnetic field. The value of this index determines the shape of the constraints for axion masses at which conversion is resonantly enhanced. For $k_H^{-1} < k_L^{-1} \ll L$, the a - γ conversion depends only on the statistical properties of the B field as worked out in Supplemental Material [79].

For low-mass axions, one finds a flat function of k_a , as anticipated earlier,

$$\langle P_{a\gamma}(k_a) \rangle = g_{a\gamma}^2 B_0^2 L \frac{\pi(\beta-1)}{8\beta k_L}, \quad \text{if } k_a \ll k_L. \quad (3)$$

Here, $B_0^2 = \langle |\mathbf{B}(\mathbf{r})|^2 \rangle$, and we use the schematic values for the injection scale k_L collected in Table I; for the LMC, this is extracted from Ref. [87]. For the MW, our choice of k_L is somewhat larger than the typical values in Refs. [69,70,74,75], which range from $16 \text{ pc} \lesssim k_L^{-1} \lesssim 175 \text{ pc}$.

For larger k_a , fewer B -field modes are on resonance, and the average conversion efficiency declines with the same power law as the turbulence spectrum,

$$\langle P_{a\gamma}(k_a) \rangle = g_{a\gamma}^2 B_0^2 L \frac{(\beta^2-1)}{4\beta(\beta+2)k_L} \left(\frac{k_L}{k_a} \right)^\beta, \quad \text{if } k_L \ll k_a \ll k_H. \quad (4)$$

For yet larger k_a , the now-familiar k_a^{-2} scaling sets in,

$$\langle P_{a\gamma}(k_a) \rangle = \frac{g_{a\gamma}^2 B_0^2}{3} \frac{1}{k_a^2}, \quad \text{if } k_H \ll k_a, \quad (5)$$

independent of any length scale associated with the B -field configuration.

LMC magnetic field—The SN 1987A host galaxy also contributes to conversion. The coherent LMC field was measured with polarized synchrotron emission at 1.4 GHz [71]. Fermi-LAT γ -ray detections imply, without assuming equipartition, $B \lesssim 7 \mu\text{G}$ after summing the plane-of-the-sky and LoS components, consistent with $2 \mu\text{G}$ expected from equipartition [71]. Slightly smaller values of $\sim 1 \mu\text{G}$ derive from Faraday rotation [72]. A larger field of $\sim 11 \mu\text{G}$ exists along two filaments, but the LoS to SN 1987A does not intercept these structures.

The small-scale LMC field is studied in Ref. [87], where Faraday rotation from 250 different sources is isolated through the structure-function method with higher-order stencils. Assuming isotropic Gaussian random fields and log-normal thermal electron densities, the magnetic energy is found to follow a power law consistent with Kolmogorov shape. The general properties that we use here are collected in Table I. Finally, the LoS depth to SN 1987A obtained from pulsar dispersion measures is $L \sim 800 \text{ pc}$ for the region RII in Table 5 of Ref. [73], where Sanduleak was located. A larger value of several kpc was gathered from red clump star observations [88], a value sometimes used in cosmic-ray studies [89]. To be conservative, we use the smaller value that we show in Table I.

SN 1987A as an axion source—Axions are emitted by the proto-neutron star (PNS) that has formed in the core after its collapse. The main relevant PNS parameters are its internal temperature (which depends on the assumed equation of state), baryon density, total mass, and cooling speed. These properties cannot be inferred from SN 1987A observations, and the neutrino signal is also not very constraining [45]. Among a suite of PNS models [31], we select their Cold Model, which produces a conservative axion emission. It has an average $T \simeq 30 \text{ MeV}$ and final baryonic mass of $1.35 M_\odot$.

For production, we assume either Primakoff scattering or nucleon bremsstrahlung. Inspired by the Kim-Shifman-Vainshtein-Zakharov model, we assume that axions couple only to protons with a structure $(g_{ap}/2m_p) \bar{p} \gamma^\mu \gamma^5 p \partial_\mu a$. We ignore pion production, $\pi^- p \rightarrow na$, because of the very uncertain π^- abundance [31]. We may then use directly the time-integrated spectral fluxes that were provided in the following form [31]:

$$\frac{dN_a}{dE_a} = N_a \frac{(\alpha+1)^{\alpha+1}}{\Gamma(\alpha+1)} \frac{E_a^\alpha}{E_{av}^{\alpha+1}} \exp \left[-\frac{(1+\alpha)E_a}{E_{av}} \right], \quad (6)$$

where N_a is the total number of emitted axions, E_{av} their average energy, $\alpha = (2 - \langle E^2 \rangle / E_{av}^2) / (\langle E^2 \rangle / E_{av}^2 - 1)$ the pinching parameter, and Γ the Gamma function

TABLE II. Axion flux parameters for Eq. (6), inspired by the Cold Model of Ref. [31], with $g = g_{a\gamma}\text{GeV}$ for Primakoff production and $g = g_{ap}$ for bremsstrahlung. $\mathcal{N}_a$ is the number of axions passing through the SMM effective area of 90 cm^2 in the 25–100 MeV channel. $P_{95}^{a\gamma}$ is the SMM limit (95% CL) on the $a\text{-}\gamma$ conversion probability.

Source reaction	$\mathcal{N}_a$ [$10^{73}g^2$]	E_{av} [MeV]	α	$\mathcal{N}_a$ [$10^{27}g^2$]	$g^2 P_{95}^{a\gamma}$ [10^{-28}]
Primakoff ($\gamma p \rightarrow pa$)	1.6	86	2.1	2.8	71
Brems ($np \rightarrow npa$)	34	59	1.6	69	2.9

(Table II). Despite a different microphysical treatment, the flux parameters of the Bari group for the Cold Model [90] agree reasonably well with ours [31].

Fluctuations of the conversion probability—The stochastic nature of turbulent fields causes fluctuations of $P_{a\gamma}$, depending on LoS and also on axion energy through k_a , as detailed in Supplemental Material [79]. This effect is particularly pronounced for massless axions ($k_a = 0$). Assuming Gaussian distributions, $\mathcal{A}_{a\gamma}$ in Eq. (1) follows an independent Gaussian for each photon polarization, implying an exponential distribution for $P_{a\gamma}(k_a)|_{k_a=0}$. The energy-averaged version $P_{a\gamma}|_{m_a=0}$ also follows an exponential distribution; therefore, it is affected by large LoS-dependent fluctuations. The physical origin is clear: at all energies, axions probe in essence the power at a single wave number ($k_a = 0$), which has large chances of fluctuating.

For nonvanishing m_a , implying a nonvanishing and energy-dependent k_a , these fluctuations diminish because resonant conversion happens for all wave numbers spanned by the energy spectrum—it is unlikely that, for a given LoS, all modes are simultaneously suppressed. Fluctuations are suppressed when $k_a \gtrsim 2\pi/L$, a regime in which the conversion probability is distributed as $\propto P_{a\gamma}^{\kappa-1} e^{-\kappa P_{a\gamma}/\langle P_{a\gamma} \rangle}$. For the large shape parameter κ , it features a much smaller dispersion compared to the exponential distribution, which itself corresponds to $\kappa = 1$. In the heavy-mass regime ($k_a \gtrsim k_H$), there is no resonance, and $P_{a\gamma}$ is once more exponentially distributed. The intermittency of turbulence, which breaks its Gaussianity, would likely increase the chances of large $P_{a\gamma}$ values, but for deriving constraints, we stick to the conservative Gaussian assumption.

We have validated our estimates with a Monte Carlo simulation of axion conversion in a synthetic turbulent Gaussian field. We have computed $P_{a\gamma}$ from 10^4 realizations for a grid of k_a values and convoluted these probabilities with the thermal spectrum of Eq. (6). The results are shown in Fig. 2. Our simulations confirm the large LoS-dependent fluctuations of $P_{a\gamma}$ for $k_a < 2\pi/L$ and $k_a \gtrsim k_H$. On the other hand, in the intermediate regime, fluctuations are indeed diminished.

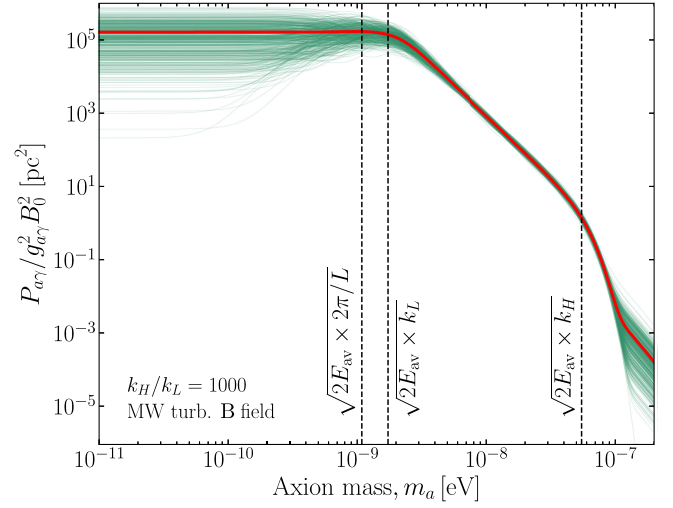


FIG. 2. Representative Monte Carlo results for $P_{a\gamma}$, the $a\text{-}\gamma$ conversion probability, for 500 of our 10^4 realizations, using the bremsstrahlung spectral axion flux. The red line is $\langle P_{a\gamma} \rangle$. Vertical lines indicate the characteristic axion masses at which a typical k_a (evaluated at $E = E_{av}$) is comparable to the inverse of the length, L , of the conversion region, the injection scale k_L , or the dissipation scale of turbulence k_H .

Results—The SMM satellite did not see any excess γ rays coincident with the neutrino burst, tagged by the time of the first event in the IMB detector, which had a reliable clock. Following Ref. [31], we use the 25–100 MeV channel with an effective area of 90 cm^2 . In Table II, we show the number of axions, $\mathcal{N}_a$, passing through this area and with energies in this range. Within three time bins (length 2.048 s each) encompassing the axion signal that lasted for less than 5 s, SMM counted $N_{\text{obs}} = 60$ events, with an expected background of $B = 54$ [31], consistent with no additional events. Assuming a Poisson distribution for N_{obs} , these values imply a constraint on a possible signal $S < 20.1$ (95% CL) [31]. One expects $S = P_{a\gamma} \mathcal{N}_a$ signal events, leading to the shown limits on $P_{a\gamma}$. In the massless limit, one finds from the coherent MW field $P_{a\gamma} = 2.1 \times 10^{-3} (g_{a\gamma}\text{GeV})^2$ (Table I), implying $|g_{ap} \times g_{a\gamma}| < 3.7 \times 10^{-24}\text{ GeV}^{-1}$ (95% CL). For larger masses, $P_{a\gamma}(k_a)$ depends on energy through $k_a = m_a^2/2E$ so that one needs to convolve the spectrum with $P_{a\gamma}(k_a)$.

For the turbulent fields, we use the parameters specified in Table I and $k_H/k_L = 1000$. We stress that, while we have observational information on k_L , the value of k_H and then of the ratio k_H/k_L cannot be measured. As we argue in Supplemental Material [79], $k_H = 1000k_L$ is a conservative choice. If this ratio were larger, the resonant enhancement would extend to larger axion masses, as the value of m_a at which the transition between the intermediate regime and the heavy regime occurs scales as $\sim \sqrt{k_H/k_L}$.

In this case, the expected number of events S is itself a random variable, which we model as a Gamma distribution

$\Gamma(\kappa, \langle S \rangle)$ and which then has to be convolved with the Poisson distribution. We use our Monte Carlo simulations to extract, for each value of m_a , the value of the shape parameter κ , using the 25–100 MeV channel of SMM. Therefore, in this case, the expected number of SMM signal events depends also on m_a . As an example, for $m_a = 3$ neV and in the MW case, we find $\bar{S} < 24.1$ (95% CL). Using the value of $\langle P_{a\gamma} \rangle$ given by our simulations, which matches our analytical predictions, this translates into $|g_{ap} \times g_{a\gamma}| < 6.4 \times 10^{-24} \text{ GeV}^{-1}$. Doing so for each m_a value leads to the constraints shown in Fig. 1.

The large fluctuations of $\bar{S}$ as a function of B -field configuration weaken the turbulent bounds considerably for small m_a . While in the massless limit, the MW turbulent $\langle P_{a\gamma} \rangle$ is larger than the coherent one, the turbulent constraint is $\bar{S} < 158.8$ (95% CL), and the bound on $|g_{ap} \times g_{a\gamma}|$ a factor of 2 weaker. For larger m_a , on the other hand, due to the resonant enhancement and the smaller fluctuations in the turbulent case, the turbulent bounds vastly supersede the coherent ones. Moreover, for a larger CL, the bounds will diminish much more for massless axions because of the larger $P_{a\gamma}$ fluctuations. This issue is further discussed in Supplemental Material [79]. For both coherent and turbulent fields, the strongest constraints derive from the MW, as shown in Fig. 1.

In Fig. 3, we show the bounds in the g_{ap} - $g_{a\gamma}$ plane for the specific value $m_a = 3$ neV, representative of the strongest enhancement by turbulence. Our results are leading constraints for $6 \times 10^{-15} \text{ GeV}^{-1} \lesssim g_{a\gamma} \lesssim 2 \times 10^{-12} \text{ GeV}^{-1}$ and $3 \times 10^{-12} \lesssim g_{ap} \lesssim 1 \times 10^{-9}$. The bound flattens out

for $g_{ap} \lesssim g_{a\gamma} m_p$, where the flux from Primakoff production prevails over bremsstrahlung [31].

The bound on $g_{a\gamma}$ shown in Fig. 3 arises from a magnetic white dwarf with a particularly low linear polarization signal amplitude [91], whereas the bound on g_{ap} from the cooling of SN 1987A [31,46]. The product of these individual limits implies $g_{ap} \times g_{a\gamma} \lesssim 2 \times 10^{-21} \text{ GeV}^{-1}$ and as such is the main competing bound to those of Fig. 1. It is comparable to the one from conversion in the assumed but unmeasured B field of the SN 1987A progenitor. A supplemental probe on $g_{a\gamma} \times g_{aN}$ is offered by the deexcitation of sodium through axion emission during the late-stage evolution of massive stars [92], but even optimistic future projections are far less sensitive than the bounds of Fig. 1.

Discussion—The nonobservation of γ rays in coincidence with SN 1987A neutrinos has led to restrictive limits on axions. Turbulent fields offer a new target for conversions. A previous work found that the average conversion probability for massless axions is boosted by up to a factor 2 [29]. While we recover this result, the constraints are actually less restrictive because turbulent conversions in this regime are highly fluctuating and cannot confidently exclude any new parameter space.

On the other hand, we identify a novel resonant target, offered by small-scale turbulent fields, for the conversion of massive axions. In this way, we extend SN 1987A sensitivity to much larger masses (Fig. 1), surpassing previous bounds by up to 2 orders of magnitude in $g_{ap} \times g_{a\gamma}$. When axions are produced by Primakoff scattering through their photon coupling alone, corresponding constraints on $g_{a\gamma}$ exist, as shown in Supplemental Material [79].

The conversion probability $P_{a\gamma}$ in such random fields arises from the resonance with the undulating modes of the turbulent B field. For axions so light that $k_a \ll k_L$, the resonance takes place for all the B -field modes, and the average conversion probability, $\langle P_{a\gamma}(k_a) \rangle$, is maximized and independent of m_a . If the axion is so heavy that k_a^{-1} is much shorter than the turbulence scales, $\langle P_{a\gamma}(k_a) \rangle$ is suppressed as m_a^{-4} . However, we also find an intermediate regime, where only a fraction of the modes are in resonance, a regime characterized by k_a between the cutoffs of the B -field power spectrum. In this regime, $\langle P_{a\gamma}(k_a) \rangle$ is suppressed as $m_a^{-2\beta}$, where the power-law index of the turbulent B field is $\beta = 5/3$ for a Kolmogorov spectrum. The intermediate regime features small dispersion around $\langle P_{a\gamma} \rangle$ since the fluctuations are averaged out when integrating over the axion spectrum. What matters most, however, is that in this regime axions maintain resonant conversion over the entire propagation length, so the conversion probability is parametrically enhanced by a rather large factor Lk_a .

A natural future extension of these results would be to apply our formalism to other kinds of turbulent magnetic fields. For example, the turbulent fields of starburst

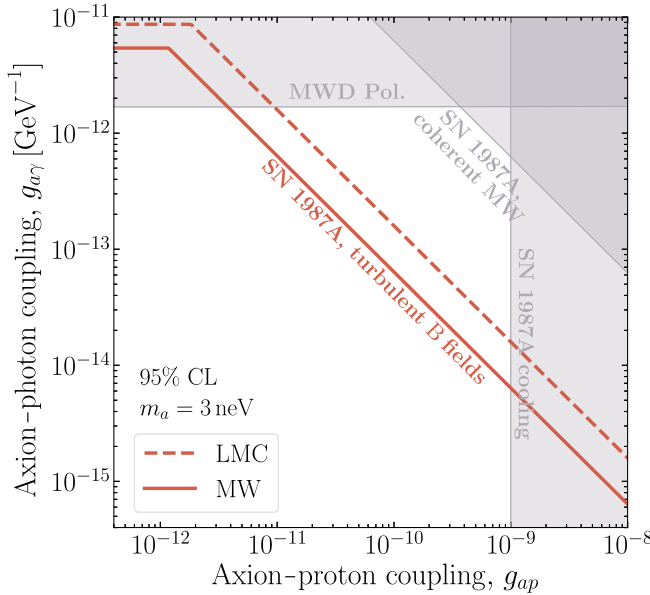


FIG. 3. Constraints in the g_{ap} - $g_{a\gamma}$ plane for fixed $m_a = 3$ neV. Existing limits include the bound from SN 1987A cooling [31,46] and from the nonobservation of axion-induced linearly polarized light from magnetic white dwarfs [91].

galaxies such as M82 and M87, recently targeted for axion searches [93–95], can potentially boost the conversion probability, extending existing bounds on, e.g., $g_{a\gamma}$ and $g_{a\gamma} \times g_{ae}$. Yet other possible targets are stellar surfaces, such as on the Sun [96], and on red supergiants (RSG) [30], the typical progenitors of core-collapse SNe. Stars can be covered by convective cells, where turbulent fields exist, so that existing constraints and projected sensitivities might be parametrically affected.

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